

NON-MINIMAL CHAOTIC INFLATION, PECCEI-QUINN PHASE TRANSITION AND NON-THERMAL LEPTOGENESIS

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ABSTRACT: We consider a phenomenological extension of the minimal supersymmetric standard model (MSSM) which incorporates non-minimal chaotic inflation, driven by a quadratic potential in conjunction with a linear term in the frame function. Inflation is followed by a Peccei-Quinn phase transition, based on renormalizable superpotential terms, which resolves the strong CP and μ problems of MSSM and provide masses lower than about 10^{12} GeV for the right-handed (RH) (s)neutrinos. Baryogenesis occurs via non-thermal leptogenesis, realized by the out-of-equilibrium decay of the RH sneutrinos, which are produced by the inflaton's decay. Confronting our scenario with the current observational data on the inflationary observables, the light neutrino masses, the baryon asymmetry of the universe and the gravitino limit on the reheat temperature, we constrain the strength of the gravitational coupling to rather large values ($\sim 45 - 2950$) and the Dirac neutrino masses to values between about 1 and 10 GeV.

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I. INTRODUCTION

There is recently a wave of interest in implementing *non-minimal chaotic inflation* (nMCI) within both a *non-supersymmetric* (SUSY) [1–8] and a SUSY [9–13] framework. The main idea is to introduce a large non-minimal coupling of the inflaton field to the curvature scalar, \mathcal{R} . After that, one can make a transformation – from the *Jordan frame* (JF) to the *Einstein* (EF) one – which flattens the potential sufficiently to support nMCI. The implementation of this mechanism within *supergravity* (SUGRA) has been greatly facilitated after the developed [10] superconformal approach to SUGRA. In particular, it is shown that the frame function can be related to a logarithmic type Kähler potential which ensures canonical kinetic terms for the scalars of the theory and incorporates an holomorphic function, F , which expresses the non-minimal coupling of the inflaton field to \mathcal{R} . Until now, the proposed models [10–13] of nMCI within SUGRA are constructed coupling quadratically the inflaton superfield with another one in the superpotential – leading thereby to a quartic potential – and adopting a quadratic term for it in F .

In this paper we propose a novel realization of nMCI within SUGRA, according to which the inflaton superfield is coupled linearly to another superfield in the superpotential of the model. As a consequence, a quadratic potential for the inflaton arises which supports nMCI, if the inflaton develops a linear coupling to \mathcal{R} . Actually, this set-up represents the SUSY implementation of the model of nMCI with $n = -1$ introduced in Ref. [7]. In contrast to earlier models [10, 14] which relied on the same superpotential term – see also Ref. [15] –, no extra shift symmetry is imposed on the Kähler potential. The resulting mass of the inflaton lies at the intermediate scale and the inflationary observables are principally similar to those of nMCI with quartic – not quadratic – potential and therefore, in excellent agreement with the current observational data [16].

The inflationary model can be nicely embedded in a mod-

est phenomenological extension of the *minimal supersymmetric standard model* (MSSM) which incorporates a resolution of the strong CP problem [17] via a *Peccei-Quinn* (PQ) symmetry. Note that there is an increasing interest [18, 19] in such models at present, since they provide us with two additional *cold dark matter* (CDM) candidates (axino and axion) beyond the lightest neutralino. In our model, a *PQ phase transition* (PQPT), tied on renormalizable [11, 20] superpotential terms, can follow nMCI generating in addition, the μ term of MSSM and intermediate masses for the *right handed* (RH) [s]neutrinos, ν_i^c [$\tilde{\nu}_i^c$]. As a consequence, the light neutrino masses can be explained through the well-known seesaw mechanism [21] provided that no large hierarchies occur in the Dirac neutrino masses. The possible formation [22] of disastrous domain walls can be avoided [23, 24] by introducing extra matter superfields without jeopardizing the gauge unification of MSSM. The appearance of a Lagrangian quartic coupling of the inflaton ensures its decay to $\tilde{\nu}_i^c$, whose the subsequent out-of-equilibrium decays can generate the *Baryon Asymmetry of the Universe* (BAU) via *non-thermal leptogenesis* (nTL) [25], consistent with the present data on neutrino data [26, 27]. Our model favors mostly quasi-degenerate ν_i^c – as in Ref. [28] – which enhances the contribution from the self-energy corrections to leptonic asymmetries, without jeopardizing the validity of the relevant perturbative results, though. The constraints arising from BAU and the gravitino (\tilde{G}) limit [29–31] on the reheat temperature can be met provided that the masses of \tilde{G} lie in the multi-TeV region.

In Sec. II we present the basic ingredients of our model, Sec. III describes the inflationary scenario, and we outline the mechanism of nTL in Sec. IV. We then restrict the model parameters in Sec. V and summarize our conclusions in Sec. VI. Throughout the text, the subscript of type χ denotes derivation *with respect to* (w.r.t) the field χ (e.g., $_{,\chi\chi} = \partial^2/\partial\chi^2$); charge conjugation is denoted by a star and brackets are, also, used by applying disjunctive correspondence.

II. MODEL DESCRIPTION

We focus on a PQ invariant extension of MSSM, which is augmented with (i) two superfields (P and \bar{P}) which are necessary for the implementation of nMCI; (ii) three superfields (S, Φ and $\bar{\Phi}$) involved in the spontaneous breaking of the PQ symmetry, $U(1)_{\text{PQ}}$ (iii) three RH neutrinos, ν_i^c , which are necessitated for the realization of the see-saw mechanism; (iv) n – to be determined below – pairs of $SU(3)_C$ triplets and antitriplets superfields, \bar{D}_a and D_a respectively, ($a = 1, \dots, n$) in order to avoid the formation of domain walls – c.f. Ref. [23, 24] – and (v) an equal number of pairs of $SU(2)_L$ doublet superfields, \bar{h}_a and h_a in order to restore gauge coupling unification at one loop – see below. Besides the superfields in the points (iv) and (v), all the others are singlets under the *Standard Model* (SM) gauge group $G_{\text{SM}} = SU(3)_C \times SU(2)_L \times U(1)_Y$. Besides the (color) anomalous $U(1)_{\text{PQ}}$, the model also possesses an anomalous R symmetry $U(1)_R$ the baryon number symmetry $U(1)_B$ and two accidental symmetries $U(1)_D$ and $U(1)_h$. The representations under G_{SM} , and the charges under the global symmetries of the various matter and Higgs superfields are listed in Table I. Note that the lepton number is not conserved in our model.

In particular, the superpotential, W , of our model can be split into four parts:

$$W = W_{\text{MSSM}} + W_{\text{DW}} + W_{\text{CPQ}} + W_{\text{NR}}, \quad (1)$$

which are analyzed in the following:

1. W_{MSSM} is the part of W which contains the usual terms – except for the μ term – of MSSM, supplemented by Yukawa interactions among the left-handed leptons and ν_i^c :

$$W_{\text{MSSM}} = h_{Dij} d_i^c Q_j H_d + h_{Uij} u_i^c Q_j H_u + h_{Eij} e_i^c L_j H_d + h_{Nij} \nu_i^c L_j H_u. \quad (2)$$

Here, the group indices have been suppressed and summation over the generation indices i and j is assumed; the i -th generation $SU(2)_L$ doublet left-handed quark and lepton superfields are denoted by Q_i and L_i respectively, and the $SU(2)_L$ singlet antiquark [antilepton] superfields by u_i^c and d_i^c [e_i^c and ν_i^c] respectively. The electroweak $SU(2)_L$ doublet Higgs superfield, which couples to the up [down] quark superfields, is denoted by H_u [H_d].

2. W_{DW} is the part of W which gives intermediate scale masses via $\langle \bar{\Phi} \rangle$ – see below – to $\bar{D}_a - D_a$ and $\bar{h}_a - h_a$. Namely,

$$W_{\text{DW}} = \lambda_{D_a} \bar{\Phi} \bar{D}_a D_a + \lambda_{h_a} \bar{\Phi} \bar{h}_a h_a. \quad (3)$$

Here, we chose a basis in the $\bar{D}_a - D_a$ and $\bar{h}_a - h_a$ space where the coupling constant matrices λ_{D_a} and λ_{h_a} are diagonal. Although these matter fields acquire intermediate scale masses after the PQ breaking, the unification of the MSSM gauge coupling constants is not disrupted at one loop. In fact, if we estimate the contribution of \bar{D}_a, D_a , and \bar{h}_a and h_a to the coefficients b_1, b_2 , and b_3 , controlling [32] the one loop evolution of the three gauge coupling constants g_1, g_2 , and g_3 , we find that the quantities $b_2 - b_1$ and $b_3 - b_2$ (which are [32] crucial for the unification of g_1, g_2 , and g_3) remain unaltered.

TABLE I: Superfield Content of the Model

Super- fields	Representations	Global Symmetries				
	under G_{SM}	R	PQ	B	D	h
Matter Fields						
L_i	$(\mathbf{1}, \mathbf{2}, -1/2)$	0	-1	0	0	0
e_i^c	$(\mathbf{1}, \mathbf{1}, 1)$	2	-1	0	0	0
ν_i^c	$(\mathbf{1}, \mathbf{1}, 0)$	2	-1	0	0	0
Q_i	$(\mathbf{3}, \mathbf{2}, 1/6)$	1	-1	$1/3$	0	0
u_i^c	$(\bar{\mathbf{3}}, \mathbf{1}, -2/3)$	1	-1	$-1/3$	0	0
d_i^c	$(\bar{\mathbf{3}}, \mathbf{1}, 1/3)$	1	-1	$-1/3$	0	0
Extra Matter Fields						
D_a	$(\mathbf{3}, \mathbf{1}, -1/3)$	1	1	0	1	0
\bar{D}_a	$(\bar{\mathbf{3}}, \mathbf{1}, 1/3)$	1	1	0	-1	0
h_a	$(\mathbf{1}, \mathbf{2}, 1/2)$	1	1	0	0	1
\bar{h}_a	$(\mathbf{1}, \mathbf{2}, -1/2)$	1	1	0	0	-1
Higgs Fields						
H_d	$(\mathbf{1}, \mathbf{2}, -1/2)$	2	2	0	0	0
H_u	$(\mathbf{1}, \mathbf{2}, 1/2)$	2	2	0	0	0
S	$(\mathbf{1}, \mathbf{1}, 0)$	4	0	0	0	0
Φ	$(\mathbf{1}, \mathbf{1}, 0)$	0	2	0	0	0
$\bar{\Phi}$	$(\mathbf{1}, \mathbf{1}, 0)$	0	-2	0	0	0
P	$(\mathbf{1}, \mathbf{1}, 0)$	6	1	0	0	0
\bar{P}	$(\mathbf{1}, \mathbf{1}, 0)$	-2	-1	0	0	0

3. W_{CPQ} is the part of W which is relevant for nMCI, the spontaneous breaking of $U(1)_{\text{PQ}}$, the decay of the inflaton and the generation of the masses of ν_i^c 's and the μ term of MSSM. It takes the form

$$W_{\text{CPQ}} = m \bar{P} P + \lambda_a S (\Phi \bar{\Phi} - M_{\text{PQ}}^2) + \lambda_{i\nu^c} \Phi \nu_i^{c2}, \quad (4)$$

where $M_{\text{PQ}} = f_a/2$ with $f_a \simeq (10^{10} - 10^{12})$ GeV being the axion decay constant which coincides with the PQ breaking scale. The parameters λ_a and f_a can be made positive by field redefinitions. From the terms in the *right hand side* (RHS) of Eq. (4) we note that the imposed symmetries disallow renormalizable terms mixing \bar{P} with some other superfields, which avoids undesirable instabilities faced in Ref. [11].

4. W_{NR} is the part of W which contains its non-renormalizable terms. Namely, we have

$$W_{\text{NR}} = \lambda_i \frac{\bar{P} S \Phi \nu_i^c}{m_P} + \lambda_P \frac{P \bar{P} \bar{\Phi} \Phi}{m_P} + \lambda_\mu \frac{\bar{\Phi}^2 H_u H_d}{m_P}, \quad (5)$$

where $m_P \simeq 2.44 \cdot 10^{18}$ GeV is the reduced Planck scale. The first term in the RHS of Eq. (5) helps accomplish sufficiently low reheating temperature and leads to the production of $\bar{\nu}_i^c$'s as dictated by nTL – see Sec. IV A. Finally, the third term provides the μ term of MSSM – see below.

To get an impression for the role that each term in the RHS of Eqs. (3), (4) and (5) play, we display the SUSY potential, V_{SUSY} , induced from the following part of W

$$W_{\text{CI}} = W_{\text{CPQ}} + W_{\text{DW}}, \quad (6)$$

which turns out to be

$$\begin{aligned}
V_{\text{SUSY}} = & m^2 (|P|^2 + |\bar{P}|^2) + |\lambda_{i\nu^c} \tilde{\nu}_i^c + \lambda_a S \bar{\Phi}|^2 \\
& + 4\lambda_{i\nu^c}^2 |\tilde{\nu}_i^c \Phi|^2 + \lambda_a^2 |\bar{\Phi} \Phi - M_{\text{PQ}}^2|^2 \\
& + |\lambda_a S \Phi + \lambda_{D_a} \bar{D}_a D_a + \lambda_{h_a} \bar{h}_a h_a|^2 \\
& + \lambda_{D_a}^2 (|\bar{D}_a|^2 + |D_a|^2) |\bar{\Phi}|^2 \\
& + \lambda_{h_a}^2 (|\bar{h}_a|^2 + |h_a|^2) |\bar{\Phi}|^2,
\end{aligned} \quad (7a)$$

where the complex scalar components of the superfields $P, \bar{P}, S, \bar{\Phi}, \Phi, \bar{D}_a, D_a, \bar{h}_a$, and h_a are denoted by the same symbol as the corresponding superfields. From Eq. (6) and assuming [20] canonical Kähler potential for the hidden sector fields, we can also derive the soft SUSY-breaking part of the inflationary potential which reads:

$$\begin{aligned}
V_{\text{soft}} = & m_{\phi^\alpha}^2 \phi^\alpha \phi_\alpha^* + \left(m B P \bar{P} - a_T \lambda_a S M_{\text{PQ}}^2 \right. \\
& + \lambda_{D_a} A_{D_a} \bar{\Phi} \bar{D}_a D_a + \lambda_{h_a} A_{h_a} \bar{\Phi} \bar{h}_a h_a \\
& \left. + \lambda_a A_a S \Phi \bar{\Phi} + \lambda_{i\nu^c} A_{i\nu^c} \Phi \nu_i^c + \text{h.c.} \right)
\end{aligned} \quad (7b)$$

where m_{ϕ^α} , with

$$\phi^\alpha = P, \bar{P}, S, \bar{\Phi}, \Phi, \tilde{\nu}_i^c, \bar{D}_{ka}, D_{ka}, \bar{h}_{la}, h_{la} \quad (8)$$

$A_a, A_{i\nu^c}, A_{D_a}, A_{h_a}, B$ and a_T are soft SUSY-breaking mass parameters of order 1 TeV. From the potential in Eqs. (7a) and (7b), we find that the SUSY vacuum lies at

$$\langle P \rangle = \langle \bar{P} \rangle = \langle \tilde{\nu}_i^c \rangle = 0, \quad (9a)$$

$$\langle D_{ka} \rangle = \langle \bar{D}_{ka} \rangle = \langle h_{la} \rangle = \langle \bar{h}_{la} \rangle = 0, \quad (9b)$$

and

$$\langle S \rangle = \frac{|A_a| + |a_T|}{2\lambda_a}, \quad |\langle \phi_\Phi \rangle| = 2|\langle \Phi \rangle| = 2|\langle \bar{\Phi} \rangle| = f_a, \quad (9c)$$

where the resulting $\langle S \rangle$ is of the order of TeV – cf. Ref. [20] – and we have introduced the canonically normalized scalar field $\phi_\Phi = 2\Phi = 2\bar{\Phi}$. Also, we use the subscripts $k = 1, 2, 3$ and $l = 1, 2$ to denote the components of D_a, \bar{D}_a and h_a, \bar{h}_a , respectively. Note that, since the sum of the arguments of $\langle \bar{\Phi} \rangle, \langle \Phi \rangle$ must be 0, $\bar{\Phi}$ and Φ can be brought to the real axis by an appropriate PQ transformation. After the spontaneous breaking of $U(1)_{\text{PQ}}$, the third term in Eq. (4) generates intermediate scale masses, $M_{i\nu^c}$ for the ν_i^c 's and, thus, seesaw masses [21] for the light neutrinos – see Sec. IV. The third term in the RHS of Eq. (5) leads to the μ term of MSSM, with $|\mu| \sim \lambda_\mu |\langle \bar{\Phi} \rangle|^2 / m_{\text{P}}$, which is of the right magnitude if $|\langle \bar{\Phi} \rangle| = f_a/2 \simeq 5 \cdot 10^{11}$ GeV, $\lambda_\mu \sim (0.001 - 0.01)$. Finally, since $\langle \bar{\Phi} \Phi \rangle / m_{\text{P}} = M_{\text{PQ}}^2 / m_{\text{P}} \ll m \simeq 10^{16}$ GeV – see Sec. III B and V B 1 – the second term in the RHS of Eq. (5) has no impact on our results.

Nonetheless, W_{CI} also gives rise to a stage of nMCI within SUGRA, if it is combined with a suitable Kähler potential, K , related to the frame function, Ω_{CI} via

$$K = -3m_{\text{P}}^2 \ln(-\Omega_{\text{CI}}/3). \quad (10)$$

In JF a specific form of Ω_{CI} 's – see Ref. [10, 11] – ensures canonical kinetic terms of the fields involved and a non-minimal coupling of the inflaton to \mathcal{R} represented by an holomorphic function $F(P)$. Going from JF to EF, and expanding the EF potential, \hat{V} , along a stable direction – usually with all the fields besides inflaton placed at the origin – \hat{V} takes the simple form

$$\hat{V}_{\text{CI0}} \simeq V_{\text{SUSY}} / f(\sigma)^2, \quad (11)$$

where $\sigma = \sqrt{2}|P|$ and f can be found expanding Ω_{CI} . Vanishing of the non-inflaton fields ensures, also, the elimination of some extra kinetic terms for scalars from the auxiliary vector fields – see Ref. [9–11].

Let us emphasize here that the coupling of P to \bar{P} is crucial in order to obtain the simple form of \hat{V}_{CI0} in Eq. (11), since only terms including derivatives of W_{CI} w.r.t \bar{P} survive in the EF SUGRA potential – see Sec. III A. This fact ensures the appearance of just one dominant power of σ in the numerator of the SUGRA scalar potential. Such a construction is not possible, e.g., for a superpotential term of the form mP^2 . Applying the strategy, described above Eq. (11) in our case, we can observe that along the direction

$$\begin{aligned}
\theta = \bar{P} = S = \bar{\Phi} = \Phi = \tilde{\nu}_i^c = \\
\bar{D}_{ka} = D_{ka} = \bar{h}_{la} = h_{la} = 0,
\end{aligned} \quad (12a)$$

with $\theta = \arg P$, V_{SUSY} in Eq. (7a) becomes

$$V_{\text{SUSY}} = \frac{1}{2} m^2 \sigma^2 + \lambda_a^2 M_{\text{PQ}}^4. \quad (12b)$$

Clearly, for $\sigma \gg f_a$, V_{SUSY} tends to a quadratic potential which can be flattened, according to Eq. (11), if f is mainly proportional to σ , i.e., if F is a linear function of P with a sizable coupling constant $c_{\mathcal{R}}$. Therefore, we are led to adopt the following frame function:

$$\Omega_{\text{CI}} = -3 + \frac{\phi^\alpha \phi_\alpha^*}{m_{\text{P}}^2} - \frac{k_{\bar{P}}}{m_{\text{P}}^4} |\bar{P}|^4 - (F(P) + F^*(P^*)), \quad (13a)$$

with ϕ^α 's defined in Eq. (8) and the non-minimal gravitational coupling

$$F = 3c_{\mathcal{R}} P / \sqrt{2} m_{\text{P}}, \quad (13b)$$

which breaks explicitly the imposed R and PQ symmetries during nMCI. In Eq. (13a) the coefficients $k_{\bar{P}}$ and $c_{\mathcal{R}}$, for simplicity, are taken real. We remark that we add the third term in the RHS of Eq. (13a) to cure the tachyonic mass problem encountered in similar models [9, 10] – see Sec. III A.

For $P \ll m_{\text{P}}$, we can show – see Sec. IV A – that an instability occurs in the PQ system which can drive a PQPT which leads to the v.e.s in Eq. (9c). Also, at the SUSY vacuum the explicit breaking of $U(1)_R \times U(1)_{\text{PQ}}$ through Eq. (13b) switches off – see Eq. (9a). A closer look, however, reveals that instanton and soft SUSY breaking effects explicitly break $U(1)_R \times U(1)_{\text{PQ}}$ to $\mathbb{Z}_2 \times \mathbb{Z}_{2(n-6)}$, as can be deduced from the solutions of the system

$$4r = 0 \pmod{2\pi} \quad \text{and} \quad 2(n-6)p - 12r = 0 \pmod{2\pi}, \quad (14)$$

where r and p are the phases of a $U(1)_R$ and $U(1)_{PQ}$ rotation respectively. Here, we take into account that the R charge of W and, thus, of all the soft SUSY breaking term is 4 and that the sum of the R [PQ] charges of the $SU(3)_C$ triplets and antitriplets is -12 [$2(n-6)$]. Note that no loop-induced PQ-violating term – as this appearing in the first paper of Ref. [13] – is detected in our case. It is then important to ensure that $\mathbb{Z}_2 \times \mathbb{Z}_{2(n-6)}$ is not spontaneously broken by $\langle \Phi \rangle$ and $\langle \bar{\Phi} \rangle$, since otherwise cosmologically disastrous domain walls are produced [22] during PQPT. This goal can be accomplished by adjusting conveniently the number n of $\bar{D}_a - D_a$ and $\bar{h}_a - h_a$ – see Table I. Indeed, when $n = 5$ or 7 we obtain $2p = 0 \pmod{2\pi}$ and therefore, $\mathbb{Z}_2 \times \mathbb{Z}_{2(n-6)}$ is not spontaneously broken by $\langle \Phi \rangle$ and $\langle \bar{\Phi} \rangle$. The residual unbroken \mathbb{Z}_2 subgroup of $U(1)_{PQ}$ can be identified with the usual matter parity of MSSM – see Table I – which prevents the rapid proton decay and ensures the stability of the *lightest SUSY particle* (LSP).

III. THE INFLATIONARY EPOCH

In Sec. III A we describe the salient features of our inflationary model and in Sec. III B we extract the inflationary observables.

A. STRUCTURE OF THE INFLATIONARY POTENTIAL

The EF F-term (tree level) SUGRA scalar potential, $\hat{V}_{\text{CI}0}$, of our model is obtained from W_{CI} in Eq. (6) and K in Eqs. (10) and (13a) by applying [9]

$$\hat{V}_{\text{CI}0} = e^{K/m_{\text{P}}^2} \left(K^{\alpha\bar{\beta}} F_{\alpha} F_{\bar{\beta}} - 3 \frac{|W_{\text{CI}}|^2}{m_{\text{P}}^2} \right), \quad (15a)$$

with

$$K_{\alpha\bar{\beta}} = K_{,\phi^{\alpha}\phi^{*\bar{\beta}}}, \quad K^{\bar{\beta}\alpha} K_{\alpha\bar{\gamma}} = \delta_{\bar{\gamma}}^{\bar{\beta}} \quad (15b)$$

and

$$F_{\alpha} = W_{\text{CI},\phi^{\alpha}} + K_{,\phi^{\alpha}} W_{\text{CI}}/m_{\text{P}}^2, \quad (15c)$$

where the ϕ^{α} 's are given in Eq. (8). From the resulting $\hat{V}_{\text{CI}0}$, we can deduce that along the field directions in Eq. (12a),

$$\hat{V}_{\text{CI}0} = \frac{m^2 m_{\text{P}}^2 x_{\sigma}^2 + 4\lambda_a^2 M_{\text{PQ}}^4/m_{\text{P}}^4}{2f^2} \simeq \frac{m^2 m_{\text{P}}^2 x_{\sigma}^2}{2f^2}, \quad (16a)$$

where $x_{\sigma} = \sigma/m_{\text{P}}$ and, according to the general recipe [9–11], the function

$$f = 1 + c_{\mathcal{R}} x_{\sigma} - x_{\sigma}^2/6 \quad (16b)$$

expresses the non-minimal coupling of σ to \mathcal{R} in JF. From Eq. (16a), we can verify that for $c_{\mathcal{R}} \gg 1$ and $x_{\sigma} \ll \sqrt{6}$, \hat{V}_{CI} develops a plateau since $M_{\text{PQ}} \ll m_{\text{P}}$ – see Sec. III B. Along the trajectory in Eq. (12a), we can estimate the constant potential energy density

$$\hat{V}_{\text{CI}0} = \frac{m^2 \sigma^2}{2f^2} \simeq \frac{m^2 m_{\text{P}}^2}{2c_{\mathcal{R}}^2}, \quad (17a)$$

and the corresponding Hubble parameter

$$\hat{H}_{\text{CI}} = \frac{\hat{V}_{\text{CI}0}^{1/2}}{\sqrt{3}m_{\text{P}}} \simeq \frac{m}{\sqrt{6}c_{\mathcal{R}}}. \quad (17b)$$

In order to check the stability of the direction in Eq. (12a) w.r.t the fluctuations of the various fields, we expand them in real and imaginary parts according to the prescription

$$P = \frac{\sigma e^{i\theta}}{\sqrt{2}} \quad \text{and} \quad X = \frac{\chi_1 + i\chi_2}{\sqrt{2}}, \quad (18a)$$

where

$$X = \bar{P}, S, \bar{\Phi}, \Phi, \tilde{\nu}_i^c, \bar{D}_{ka}, D_{ka}, \bar{h}_{la}, h_{la} \quad (18b)$$

and

$$\chi = \bar{p}, s, \bar{\phi}, \phi, \nu_i, \bar{D}_{ka}, D_{ka}, \bar{h}_{la}, h_{la}, \quad (18c)$$

respectively. Along the trajectory in Eq. (12a) we find

$$(K_{\alpha\bar{\beta}}) = \text{diag} \left(J^2, \underbrace{1/f, \dots, 1/f}_{7+10n \text{ elements}} \right), \quad (19a)$$

where

$$J = \sqrt{\frac{1}{f} + \frac{3}{2} m_{\text{P}}^2 \left(\frac{f_{,\sigma}}{f} \right)^2} = \frac{\sqrt{2+3c_{\mathcal{R}}^2}}{\sqrt{2}f} \simeq \sqrt{\frac{3}{2}} \frac{1}{x_{\sigma}}. \quad (19b)$$

Consequently, we can introduce the EF canonically normalized fields, $\hat{\sigma}$, $\hat{\theta}$ and $\hat{\chi}$, as follows – cf. Ref. [9–12]:

$$K_{\alpha\bar{\beta}} \dot{\phi}^{\alpha} \dot{\phi}^{*\bar{\beta}} = \frac{1}{2} \left(\dot{\hat{\sigma}}^2 + \dot{\hat{\theta}}^2 \right) + \frac{1}{2} \sum_{\chi} \left(\dot{\hat{\chi}}_1^2 + \dot{\hat{\chi}}_2^2 \right), \quad (20)$$

where the dot denotes derivation w.r.t the JF cosmic time, t and the hatted fields are defined as follows

$$\frac{d\hat{\sigma}}{d\sigma} = J, \quad \hat{\theta} \simeq J\sigma\theta \quad \text{and} \quad \hat{\chi} \simeq \frac{\chi}{\sqrt{f}}. \quad (21)$$

Note that $\hat{\theta} \simeq J\sigma\theta$ since $J\sigma \simeq \sqrt{3/2}m_{\text{P}}$ – see Eq. (19b). On the other hand, we can show that during a stage of slow-roll nMCI, $\dot{\hat{\chi}} \simeq \dot{\chi}/\sqrt{f}$ since the quantity $\dot{f}/2f^{3/2}\chi$, involved in relating $\dot{\hat{\chi}}$ to $\dot{\chi}$, turns out to be negligibly small compared with $\dot{\hat{\chi}}$. Indeed, the $\hat{\chi}$'s acquire effective masses $m_{\hat{\chi}} \gg \hat{H}_{\text{CI}}$ – see below – and therefore enter a phase of oscillations about $\hat{\chi} = 0$ with reducing amplitude. Neglecting the oscillating part of the relevant solutions, we find

$$\chi \simeq \hat{\chi}_0 \sqrt{f} e^{-2\hat{N}/3} \quad \text{and} \quad \dot{\hat{\chi}} \simeq -2\hat{\chi}_0 \sqrt{f} \hat{H}_{\text{CI}} \hat{\eta}_{\hat{\chi}} e^{-2\hat{N}/3}, \quad (22)$$

where $\hat{\chi}_0$ represents the initial amplitude of the oscillations, $\hat{\eta}_{\hat{\chi}} = m_{\hat{\chi}}^2/3\hat{H}_{\text{CI}}$ and we assume $\hat{\chi}(t=0) = 0$. Taking into account the approximate expressions for $\dot{\sigma}$ and the slow-roll parameter $\hat{\epsilon}$, which are displayed in Sec. III B, we find

$$-\frac{\dot{f}}{2f^{3/2}}\chi = \frac{c_{\mathcal{R}}\hat{\epsilon}\hat{H}_{\text{CI}}^2}{m_{\hat{\chi}}^2}\dot{\hat{\chi}} \ll \dot{\hat{\chi}}. \quad (23)$$

TABLE II: The mass spectrum of the model during nMCI

Fields	Eigenstates	Mass Squared
Bosons		
1 real scalar	$\hat{\theta}$	$c_{\mathcal{R}} m^2 x_{\sigma} / f^3 J^2 \simeq 4H_{\text{CI}}^2$
2 real scalars	\hat{p}_1, \hat{p}_2	$m^2 (2k_{\bar{P}} c_{\mathcal{R}} x_{\sigma}^3 + (-c_{\mathcal{R}} x_{\sigma} + (6k_{\bar{P}} - 1)c_{\mathcal{R}}^2 x_{\sigma}^2) / 2J^2 f^2) / f^2$
2 real scalars	\hat{s}_1, \hat{s}_2	$(2\lambda_a M_{\text{PQ}}^4 / m_{\text{P}}^2 + m^2 x_{\sigma}^2) / 3f^2$
6 real scalars	$\hat{\nu}_{1i}, \hat{\nu}_{2i}$	$(2\lambda_a M_{\text{PQ}}^4 / m_{\text{P}}^2 + m^2 x_{\sigma}^2) / f^2$
4 real scalars	$\frac{\hat{\phi}_1 \pm \hat{\phi}_1}{\sqrt{2}}, \frac{\hat{\phi}_2 \pm \hat{\phi}_2}{\sqrt{2}}$	$(m^2 x_{\sigma}^2 / 3 \pm \lambda_a^2 M_{\text{PQ}}^2 f) / f^2$
6n real scalars	$\hat{D}_{1ka}, \hat{D}_{2ka}$	$(2\lambda_a M_{\text{PQ}}^4 / m_{\text{P}}^2 + m^2 x_{\sigma}^2) / f^2$
6n real scalars	$\hat{\bar{D}}_{1ka}, \hat{\bar{D}}_{2ka}$	$(2\lambda_a M_{\text{PQ}}^4 / m_{\text{P}}^2 + m^2 x_{\sigma}^2) / f^2$
4n real scalars	$\hat{h}_{1la}, \hat{h}_{2la}$	$(2\lambda_a M_{\text{PQ}}^4 / m_{\text{P}}^2 + m^2 x_{\sigma}^2) / f^2$
4n real scalars	$\hat{\bar{h}}_{1la}, \hat{\bar{h}}_{2la}$	$(2\lambda_a M_{\text{PQ}}^4 / m_{\text{P}}^2 + m^2 x_{\sigma}^2) / f^2$
Fermions		
2 Weyl spinors	$\frac{\psi_{\text{P}} \pm \psi_{\text{P}}}{\sqrt{2}}$	$m^2 (6 + x_{\sigma}^2) / 36f^2$

The masses that the various scalars acquire during nMCI are presented in Table II. To this end, we expand \hat{V}_{CIO} in Eq. (15a) to quadratic order in the fluctuations around the direction of Eq. (12a). As we observe from the relevant eigenvalues of the mass-squared matrices, no instability – as the one found in Ref. [11] – arises in the spectrum. In particular, it is evident that $k_{\bar{P}} \gtrsim 1$ assists us to achieve positivity of the mass-squared associated with the scalars $\hat{p}_{1,2}, m_{\bar{P}}^2$ – in accordance with the results of Ref. [9, 10]. It is remarkable that mass-squared corresponding to $\hat{\nu}_i, D_{ka}, \bar{D}_{ka}, h_{la}, \bar{h}_{la}$ are independent of the relevant superpotential couplings $\lambda_{i\nu^c}, \lambda_{Da}$ and λ_{ha} . We have also numerically verified that the various masses remain greater than \hat{H}_{CI} during the last 50 – 60 e-foldings of nMCI, and so any inflationary perturbations of the fields other than the inflaton are safely eliminated.

In Table II we also present the masses squared of chiral fermions of the model along the direction of Eq. (12a). Inserting these masses into the well-known Coleman-Weinberg formula [33], we can find the one-loop radiative corrections, V_{rc} , in our model which can be written as

$$\begin{aligned}
V_{\text{rc}} = & \frac{1}{64\pi^2} \left(m_{\hat{\theta}}^4 \ln \frac{m_{\hat{\theta}}^2}{\Lambda^2} + 2m_{\hat{p}}^4 \ln \frac{m_{\hat{p}}^2}{\Lambda^2} + 2m_{\hat{s}}^4 \ln \frac{m_{\hat{s}}^2}{\Lambda^2} \right. \\
& + 6m_{\hat{\nu}}^4 \ln \frac{m_{\hat{\nu}}^2}{\Lambda^2} + 2m_{\hat{\phi}_+}^4 \ln \frac{m_{\hat{\phi}_+}^2}{\Lambda^2} + 2m_{\hat{\phi}_-}^4 \ln \frac{m_{\hat{\phi}_-}^2}{\Lambda^2} \\
& \left. + 20n m_{\hat{D}}^4 \ln \frac{m_{\hat{D}}^2}{\Lambda^2} - 4m_{\hat{\psi}_{\pm}}^4 \ln \frac{m_{\hat{\psi}_{\pm}}^2}{\Lambda^2} \right), \quad (24)
\end{aligned}$$

where $\Lambda = m_{\text{P}}/c_{\mathcal{R}}$ is the cutoff scale of the effective theory – see Sec. V A – and the involved above masses squared $m_{\hat{\theta}}^2, m_{\hat{p}}^2, m_{\hat{s}}^2, m_{\hat{\nu}}^2, m_{\hat{\phi}_{\pm}}^2, m_{\hat{D}}^2$ and $m_{\hat{\psi}_{\pm}}^2$ are equal to the ones listed in the third column of Table II from top to the bottom.

Note that the masses squared of all the extra matter fields are equal to $m_{\hat{D}}^2$. Based on the one-loop corrected EF potential

$$\hat{V}_{\text{CI}} = \hat{V}_{\text{CIO}} + V_{\text{rc}}, \quad (25)$$

we can proceed to the analysis of nMCI in EF, employing the standard slow-roll approximation [34]. It can be shown [35] that the results calculated this way are the same as if we had calculated them using the non-minimally coupled scalar field in JF. As expected and verified numerically, V_{rc} does not affect the inflationary dynamics and predictions, in the major part of the allowed parameter space – see Sec. V B 1 – since the inflationary path already possesses a slope at the classical level – see below.

B. THE INFLATIONARY OBSERVABLES

According to our analysis above, the universe undergoes a period of slow-roll nMCI, which is determined by the condition – see e.g. Ref. [34]:

$$\max\{\hat{\epsilon}(\sigma), |\hat{\eta}(\sigma)|\} \leq 1,$$

where

$$\hat{\epsilon} = \frac{m_{\text{P}}^2}{2} \left(\frac{\hat{V}_{\text{CI},\sigma}}{\hat{V}_{\text{CI}}} \right)^2 = \frac{m_{\text{P}}^2}{2J^2} \left(\frac{\hat{V}_{\text{CI},\sigma}}{\hat{V}_{\text{CI}}} \right)^2 \simeq \frac{4m_{\text{P}}^2}{3c_{\mathcal{R}}^2 \sigma^2}, \quad (26a)$$

and

$$\begin{aligned}
\hat{\eta} &= m_{\text{P}}^2 \frac{\hat{V}_{\text{CI},\sigma\sigma}}{\hat{V}_{\text{CI}}} = \frac{m_{\text{P}}^2}{J^2} \left(\frac{\hat{V}_{\text{CI},\sigma\sigma}}{\hat{V}_{\text{CI}}} - \frac{\hat{V}_{\text{CI},\sigma}}{\hat{V}_{\text{CI}}} \frac{J_{,\sigma}}{J} \right) \\
&\simeq -4m_{\text{P}}/3c_{\mathcal{R}}\sigma. \quad (26b)
\end{aligned}$$

Here, we employ Eqs. (17a) and (19b) and the following approximate relations:

$$\hat{V}_{\text{CI},\sigma} \simeq \frac{m^2 m_{\text{P}}}{c_{\mathcal{R}}^3 x_{\sigma}^2} \quad \text{and} \quad \hat{V}_{\text{CI},\sigma\sigma} \simeq -\frac{2m^2}{c_{\mathcal{R}}^3 x_{\sigma}^3}. \quad (27)$$

The numerical computation reveals that nMCI terminates due to the violation of the $\hat{\epsilon}$ criterion at $\sigma = \sigma_{\text{f}}$, which is calculated to be

$$\hat{\epsilon}(\sigma_{\text{f}}) = 1 \quad \Rightarrow \quad \sigma_{\text{f}} = \frac{2}{\sqrt{3}} \frac{m_{\text{P}}}{c_{\mathcal{R}}}. \quad (28)$$

We note, in passing, that for $\sigma \geq \sigma_{\text{f}}$ the evolution of $\hat{\sigma}$ – or σ via Eq. (19b) – is governed by the equation of motion

$$3\hat{H}_{\text{CI}} \frac{d\hat{\sigma}}{dt} = -\hat{V}_{\text{CI},\hat{\sigma}} \Rightarrow \dot{\sigma} = -\frac{2\sqrt{2}m}{3\sqrt{3}} \sqrt{\frac{m_{\text{P}}^3}{c_{\mathcal{R}}^3 \sigma}}, \quad (29)$$

where \hat{t} is the EF cosmic time with $d\hat{t} = \sqrt{f}dt$. Using Eqs. (26a) and (29), we can derive Eq. (23).

The number of e-foldings, \hat{N}_* , that the scale $k_* = 0.002/\text{Mpc}$ suffers during nMCI can be calculated through the relation

$$\hat{N}_* = \frac{1}{m_{\text{P}}^2} \int_{\hat{\sigma}_{\text{f}}}^{\hat{\sigma}^*} d\hat{\sigma} \frac{\hat{V}_{\text{CI}}}{\hat{V}_{\text{CI},\hat{\sigma}}} = \frac{1}{m_{\text{P}}^2} \int_{\sigma_{\text{f}}}^{\sigma^*} d\sigma J^2 \frac{\hat{V}_{\text{CI}}}{\hat{V}_{\text{CI},\sigma}}, \quad (30)$$

where $\sigma_* [\hat{\sigma}_*]$ is the value of $\sigma [\hat{\sigma}]$ when k_* crosses the inflationary horizon. Given that $\sigma_f \ll \sigma_*$, we can write σ_* as a function of \hat{N}_* as follows

$$\hat{N}_* \simeq \frac{3c_{\mathcal{R}}}{4m_{\text{P}}} (\sigma_* - \sigma_f) \Rightarrow \sigma_* \simeq \frac{4\hat{N}_*}{3c_{\mathcal{R}}} m_{\text{P}}. \quad (31)$$

The power spectrum $\Delta_{\mathcal{R}}^2$ of the curvature perturbations generated by σ at the pivot scale k_* is estimated as follows

$$\Delta_{\mathcal{R}} = \frac{1}{2\pi m_{\text{P}}^2} \sqrt{\frac{\hat{V}_{\text{CI}}(\sigma_*)}{6\hat{\epsilon}(\sigma_*)}} \simeq \frac{m\hat{N}_*}{6\pi m_{\text{P}} c_{\mathcal{R}}}, \quad (32)$$

where Eq. (31) is employed to derive the last equality of the relation above. Since the scalars listed in Table II are massive enough during nMCI, $\Delta_{\mathcal{R}}$ can be identified with its central observational value – see Sec. V – with almost constant \hat{N}_* . The resulting relation reveals that m is to be proportional to $c_{\mathcal{R}}$. Indeed we find

$$m = 6\pi m_{\text{P}} c_{\mathcal{R}} \Delta_{\mathcal{R}} / \hat{N}_* \Rightarrow m = 4.1 \cdot 10^{13} c_{\mathcal{R}} \text{ GeV}, \quad (33)$$

for $\hat{N}_* \simeq 55$. At the same pivot scale, we can also calculate the (scalar) spectral index, n_s , its running, a_s , and the scalar-to-tensor ratio, r , via the relations:

$$n_s = 1 - 6\hat{\epsilon}_* + 2\hat{\eta}_* \simeq 1 - 2/\hat{N}_*, \quad (34a)$$

$$\alpha_s = \frac{2}{3} (4\hat{\eta}_*^2 - (n_s - 1)^2) - 2\hat{\xi}_* \simeq -2/\hat{N}_*^2, \quad (34b)$$

$$r = 16\hat{\epsilon}_* \simeq 12/\hat{N}_*^2, \quad (34c)$$

where $\hat{\xi} = m_{\text{P}}^4 \hat{V}_{\text{CI},\hat{\sigma}} \hat{V}_{\text{CI},\hat{\sigma}\hat{\sigma}} / \hat{V}_{\text{CI}}^2 = m_{\text{P}} \sqrt{2\hat{\epsilon}} \hat{\eta}_{,\sigma} / J + 2\hat{\eta}\hat{\epsilon}$ and the variables with subscript $*$ are evaluated at $\sigma = \sigma_*$. Comparing the results of this section with the observationally favored values [16], we constrain the parameters of our model in Sec. V B 1.

IV. NON-THERMAL LEPTOGENESIS

A complete SUSY inflationary scenario should specify the transition to the radiation dominated era and also explain the origin of the observed BAU consistently with the \tilde{G} constraint. These goals can be accomplished within our set-up, as we describe in this section. Namely, the basic features of the post-inflationary evolution are exhibited in Sec. IV A and the topic of nTL in conjunction with the present neutrino data is analyzed in Sec. IV B.

A. THE GENERAL SET-UP

When nMCI is over, the inflaton continues to roll down towards the SUSY vacuum, Eqs. (9a), (9b) and (9c). Note that when $x_{\sigma} \lesssim \sqrt{3}\lambda_a M_{\text{PQ}}/m$, one scalar originating from the superfields Φ and $\bar{\Phi}$ – see Table I – acquires a negative mass squared triggering thereby the PQPT. As the inflaton continues its rolling, there is a brief stage of tachyonic preheating

[36] which does not lead to significant particle production [37]. Soon afterwards, it settles into a phase of damped oscillations about the minimum of the $\hat{V}_{\text{CI}0}$. Since no gauge symmetry is broken during nMCI, no superheavy bosons are produced and therefore no particle production via the mechanism of instant preheating [38] occurs. Also, since the inflaton cannot decay via renormalizable interactions to SM particles, effects of narrow parametric resonance [36] are also absent in our regime.

Nonetheless, the standard perturbative approach to the inflaton decay provides a very efficient decay rate. Namely, at the SUSY vacuum the fields involved acquire the v.e.s shown in Eqs. (9a), (9b) and (9c) giving rise to the mass spectrum presented in Table III. There we can show the mass, m_I , of the (canonically normalized) inflaton \hat{P} and the masses $M_{i\nu^c}$ of the RH [s]neutrinos, $\nu_i^c [\tilde{\nu}_i^c]$, which play a crucial role in our scenario of nTL. Note that since $\langle \Phi \rangle = \langle \bar{\Phi} \rangle = M_{\text{PQ}} \ll m_{\text{P}}$, $\langle \Omega \rangle \simeq -3$ and so $\langle f \rangle \simeq 1$. Therefore, apart from \hat{P} , the EF canonically normalized field are not distinguished from the JF ones at the SUSY vacuum. On the other hand, \hat{P} can be expressed as a function of P through the relation

$$\frac{\hat{P}}{P} = \langle J \rangle \quad \text{where} \quad \langle J \rangle = \sqrt{1 + 3c_{\mathcal{R}}^2/2}. \quad (35)$$

Making use of Eq. (33) we can infer that m_I is kept independent of $c_{\mathcal{R}}$ and almost constant at the level of 10^{13} GeV. Indeed,

$$m_I \simeq \frac{m}{\langle J \rangle} = \sqrt{\frac{2}{3}} \frac{m}{c_{\mathcal{R}}} \simeq 2\sqrt{3}\pi m_{\text{P}} \frac{\Delta_{\mathcal{R}}}{\hat{N}} \simeq 10^{13} \text{ GeV}, \quad (36)$$

where the WMAP7 value of $\Delta_{\mathcal{R}}$ – see Sec. V A – is employed in the last step of the relation above. In the expressions of the various eigenstates listed in Table III, we adopt the following abbreviations

$$\delta\Phi = \Phi - M_{\text{PQ}}, \quad \delta\bar{\Phi} = \bar{\Phi} - M_{\text{PQ}}, \quad (37a)$$

and

$$\psi_{\pm} = (\psi_{\bar{\Phi}} \pm \psi_{\Phi}) / \sqrt{2}, \quad (37b)$$

where ψ_x with $x = \hat{P}, P, S, \bar{\Phi}, \Phi, \bar{D}_{ka}, D_{ka}, \bar{h}_{la}$, and h_{la} denote the chiral fermions associated with the superfields $\hat{P}, P, S, \bar{\Phi}, \Phi, \bar{D}_{ka}, D_{ka}, h_{la}$, and h_{la} respectively. The eigenstates ψ_- and $\delta\Phi_-$, with

$$\delta\Phi_{\pm} = (\delta\bar{\Phi} \pm \delta\Phi) / \sqrt{2}, \quad (38)$$

contain the components of the axion supermultiplet. Namely axion [saxion] can be identified with the phase [modulus] of the complex field $\delta\Phi_-$, whereas ψ_- can be interpreted as the axino. Note that the zero masses of saxion and axino can be replaced with masses of order 1 TeV if we take into account the soft SUSY breaking masses – see discussion below Eq. (9c).

The decay of \hat{P} commences when m_I becomes larger than the expansion rate and is processed via the first coupling in

TABLE III: The mass spectrum of the model at the SUSY vacuum

Eigenstates		Eigenvalues
Scalars	Fermions	(Masses)
\hat{P}, \bar{P}	$(\psi_{\bar{P}} \pm \psi_{\bar{P}})/\sqrt{2}$	$m_I = m/\langle J \rangle$
$S, (\delta\bar{\Phi} + \delta\Phi)/\sqrt{2}$	$(\psi_S \pm \psi_+)/\sqrt{2}$	$m_{PQ} = \sqrt{2}\lambda_a M_{PQ}$
$(\delta\bar{\Phi} - \delta\Phi)/\sqrt{2}$	$(\psi_{\bar{\Phi}} - \psi_{\Phi})/\sqrt{2}$	0
$\tilde{\nu}_i^c$	ν_i^c	$M_{i\nu^c} = 2\lambda_{i\nu^c} M_{PQ}$
D_{ka}, \bar{D}_{ka}	$\psi_{D_{ka}}, \psi_{\bar{D}_{ka}}$	$m_{D_a} = \lambda_{D_a} M_{PQ}$
h_{la}, \bar{h}_{la}	$\psi_{h_{la}}, \psi_{\bar{h}_{la}}$	$m_{h_a} = \lambda_{h_a} M_{PQ}$

the RHS of Eq. (5), into S and $\tilde{\nu}_i^c$ and $S, \tilde{\nu}_i^c$ and $\delta\Phi_+$ or $\delta\Phi_-$. The relevant Lagrangian sector is

$$\mathcal{L}_{dc} = -\frac{m_I}{m_P} \lambda_i \hat{P}^* S \tilde{\nu}_i^c \left(M_{PQ} + \frac{\delta\Phi_+ - \delta\Phi_-}{\sqrt{2}} \right) + \text{h.c.} \quad (39)$$

which arises from the cross term of the F-term, corresponding to \bar{P} , of the SUSY potential derived from the superpotential terms in Eqs. (4) and (5). Note that we have no $c_{\mathcal{R}}$ -induced decay channels as in Ref. [12], since $\langle P \rangle = 0$. The interaction above gives rise to the following decay width

$$\Gamma_I = \frac{1}{8\pi} \left(\left(\frac{M_{PQ}}{m_P} \right)^2 + \frac{1}{64\pi^2} \left(\frac{m_I}{m_P} \right)^2 \right) m_I \sum_{i=1}^3 \lambda_i^2, \quad (40)$$

where we take into account that $m_I \gg m_{PQ}$ and $m_I \gg M_{j\nu^c}$. These prerequisites are safely fulfilled when λ_a and $\lambda_{i\nu^c}$ remain perturbative, i.e. $\lambda_a, \lambda_{i\nu^c} \leq \sqrt{4\pi}$ – see Table III. From the two contributions to Γ_I , the dominant one is the second one – the 3-body decay channel – originating from the two last terms of Eq. (39).

Taking also into account that the decay width of the produced $\tilde{\nu}_i^c$, $\Gamma_{i\nu^c}$, is much larger than Γ_I – see below – we can infer that the reheat temperature, T_{rh} , is exclusively determined by the \hat{P} decay and is given by [39]

$$T_{rh} = \left(\frac{72}{5\pi^2 g_*} \right)^{1/4} \sqrt{\Gamma_I m_P}, \quad (41)$$

where $g_* \simeq 232.5$ counts the effective number of the relativistic degrees of freedom at temperature T_{rh} for the (s)particle spectrum of MSSM plus the particle content of the axion supermultiplet. Although the factor before the square root of Eq. (41) differs [39] slightly from other calculations of T_{rh} – cf. Ref. [25] – the numerical result remains pretty stable and close to 10^8 GeV – see Sec. V B 1.

If $T_{rh} \ll M_{i\nu^c}$, the out-of-equilibrium condition [40] for the implementation of nTL is automatically satisfied. Subsequently, $\tilde{\nu}_i^c$ decay into \tilde{H}_u and L_i or \tilde{H}_u^* and \tilde{L}_i^* via the tree-level couplings derived from the second term of the second line of Eq. (2). Interference between tree-level and one-loop diagrams generates a lepton-number asymmetry (per ν_i^c decay) ε_i [40], when CP is not conserved in the Yukawa coupling constants h_{Nij} – see Eq. (2). The resulting lepton-number asymmetry after reheating can be partially converted

through sphaleron effects into baryon-number asymmetry. However, the required T_{rh} must be compatible with constraints for the \tilde{G} abundance, $Y_{\tilde{G}}$, at the onset of *nucleosynthesis* (BBN). In particular, the B yield can be computed as

$$Y_B = -0.35 \frac{5}{4} \frac{T_{rh}}{m_I} \sum_i \text{Br}_i \varepsilon_i \quad \text{with} \quad \text{Br}_i = \frac{\lambda_i^2}{\sum_i \lambda_i^2} \quad (42)$$

the branching ratio of \hat{P} to $\tilde{\nu}_i^c$ – see Eq. (40). In the formula above the first numerical factor (0.35) comes from the sphaleron effects, whereas the second one (5/4) is due to the slightly different calculation [39] of T_{rh} – cf. Ref. [25]. On the other hand, the \tilde{G} yield due to thermal production at the onset of BBN is estimated to be [31]:

$$Y_{\tilde{G}} \simeq 1.9 \cdot 10^{-22} T_{rh}/\text{GeV}. \quad (43)$$

where we assume that \tilde{G} is much heavier than the gauginos. Let us note that non-thermal \tilde{G} production within SUGRA is unlikely in our scenario, since these contributions are [41] usually proportional to the v.e.v of the inflaton which is zero in our case.

Both Eqs. (42) and (43) calculate the correct values of the B and \tilde{G} abundances provided that no entropy production occurs for $T < T_{rh}$ – see also Sec. V A. This fact can be easily achieved within our setting. Indeed, following the arguments of Ref. [11], one can show that the PQ system comprised of the fields S and $\delta\Phi_+$ decays via the third term in the RHS of Eq. (5) before its domination over radiation, for all relevant values of λ_i 's. Regarding the saxion, $\delta\Phi_-$, we can assume that it has mass of the order of 1 TeV, its decay mode to axions is suppressed (w.r.t the ones to gluons, higgses and higgsinos [19, 42, 43]) and the initial amplitude of its oscillations is equal to f_a . Under these circumstances, it can [42] decay before domination too, and evades [43] the constraints from the effective number of neutrinos for the f_a 's and T_{rh} 's encountered in our model. As a consequence of its relatively large decay temperature, the LSPs produced by the saxion decay are likely to be thermalized and therefore, no upper bound on the saxion abundance is [43] to be imposed. Finally, if axino is sufficiently light it can act as a CDM candidate [18, 19] with relic abundance produced predominantly thermally – due to the relatively large T_{rh} . Otherwise, it may enhance [19] non-thermally the abundance of a higgsino-like neutralino-LSP, rendering it a successful CDM candidate.

B. LEPTON-NUMBER ASYMMETRY AND NEUTRINO MASSES

As mentioned above, the decay of $\tilde{\nu}_i^c$, emerging from the \hat{P} decay, can generate a lepton asymmetry, ε_i , caused by the interference between the tree and one-loop decay diagrams, provided that a CP-violation occurs in h_{Nij} 's. The produced ε_i can be expressed in terms of the Dirac mass matrix of ν_i , m_D , defined in a basis (called ν_i^c -basis henceforth) where ν_i^c

are mass eigenstates, as follows:

$$\varepsilon_i = \frac{\sum_{j \neq i} \text{Im} \left[(m_D^\dagger m_D)_{ij}^2 \right] \left(F_S(x_{ij}, y_i, y_j) + F_V(x_{ij}) \right)}{8\pi \langle H_u \rangle^2 (m_D^\dagger m_D)_{ii}}, \quad (44a)$$

where we take $\langle H_u \rangle \simeq 174$ GeV, for large $\tan \beta$ and

$$x_{ij} = \frac{M_{j\nu^c}}{M_{i\nu^c}} \quad \text{and} \quad y_i = \frac{\Gamma_{i\nu^c}}{M_{i\nu^c}} = \frac{(m_D^\dagger m_D)_{ii}}{8\pi \langle H_u \rangle^2}. \quad (44b)$$

Also F_V and F_S represent, respectively, the contributions from vertex and self-energy diagrams which in SUSY theories read [44–46]

$$F_V(x) = -x \ln(1 + x^{-2}), \quad (44c)$$

and

$$F_S(x, y, z) = \frac{-2x(x^2 - 1)}{(x^2 - 1 - x^2 z \ln x^2 / \pi)^2 + (x^2 z - y)^2}, \quad (44d)$$

with the latter expression written as given in Ref. [46]. When

$$\Delta_{iji} \gg 1 \quad \text{and} \quad \Delta_{ijj} \gg 1 \quad \text{with} \quad \Delta_{ijk} = \frac{|x_{ij}^2 - 1|}{x_{ik} y_k}, \quad (45)$$

(no summation is applied over the repeated indices) we can simplify F_S expanding it close to $x \simeq 1$ as follows

$$F_S \simeq \frac{2x}{1 - x^2} \simeq \frac{1}{1 - x} - \frac{1}{2}. \quad (46)$$

The involved in Eq. (44a) m_D can be diagonalized if we define a basis – called weak basis henceforth – in which the lepton Yukawa couplings and the $SU(2)_L$ interactions are diagonal in the space of generations. In particular we have

$$U^\dagger m_D U^{c\dagger} = d_D = \text{diag}(m_{1D}, m_{2D}, m_{3D}), \quad (47)$$

where U and U^c are 3×3 unitary matrices which relate L_i and ν_i^c (in the ν_i^c -basis) with the ones L'_i and $\nu_i^{c'}$ in the weak basis as follows:

$$L' = LU \quad \text{and} \quad \nu^{c'} = U^c \nu^c. \quad (48)$$

Here, we write LH lepton superfields, i.e. $SU(2)_L$ doublet leptons, as row 3-vectors in family space and RH anti-lepton superfields, i.e. $SU(2)_L$ singlet anti-leptons, as column 3-vectors. Consequently, the combination $m_D^\dagger m_D$ appeared in Eq. (44a) turns out to be a function just of d_D and U^c . Namely,

$$m_D^\dagger m_D = U^{c\dagger} d_D^\dagger d_D U^c. \quad (49)$$

The connection of the leptogenesis scenario with the low energy neutrino data can be achieved through the seesaw formula, which gives the light-neutrino mass matrix m_ν in terms of m_{iD} and $M_{i\nu^c}$. Working in the ν_i^c -basis, we have

$$m_\nu = -m_D d_{\nu^c}^{-1} m_D^T, \quad (50)$$

where

$$d_{\nu^c} = \text{diag}(M_{1\nu^c}, M_{2\nu^c}, M_{3\nu^c}) \quad (51)$$

with $M_{1\nu^c} \leq M_{2\nu^c} \leq M_{3\nu^c}$ real and positive. Solving Eq. (47) w.r.t m_D and inserting the resulting expression in Eq. (50) we extract the mass matrix

$$\bar{m}_\nu = U^\dagger m_\nu U^* = -d_D U^c d_{\nu^c}^{-1} U^{cT} d_D, \quad (52)$$

which can be diagonalized by the unitary PMNS matrix satisfying

$$\bar{m}_\nu = U_\nu^* \text{diag}(m_{1\nu}, m_{2\nu}, m_{3\nu}) U_\nu^\dagger \quad (53)$$

and parameterized as follows:

$$U_\nu = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ U_{21\nu} & U_{22\nu} & s_{23}c_{13} \\ U_{31\nu} & U_{32\nu} & c_{23}c_{13} \end{pmatrix} \cdot \mathcal{P}. \quad (54)$$

Here

$$U_{21\nu} = -c_{23}s_{12} - s_{23}c_{12}s_{13}e^{i\delta}, \quad (55a)$$

$$U_{22\nu} = c_{23}c_{12} - s_{23}s_{12}s_{13}e^{i\delta}, \quad (55b)$$

$$U_{31\nu} = s_{23}s_{12} - c_{23}c_{12}s_{13}e^{i\delta}, \quad (55c)$$

$$U_{32\nu} = -s_{23}c_{12} - c_{23}s_{12}s_{13}e^{i\delta}, \quad (55d)$$

with $c_{ij} := \cos \theta_{ij}$, $s_{ij} := \sin \theta_{ij}$ and δ the CP-violating Dirac phase. The two CP-violating Majorana phases φ_1 and φ_2 are contained in the matrix

$$\mathcal{P} = \text{diag}(e^{-i\varphi_1/2}, e^{-i\varphi_2/2}, 1). \quad (56)$$

Following a bottom-up approach, along the lines of Ref. [47], we can find \bar{m}_ν via Eq. (53) using as input parameters the low energy neutrino observables, the CP violating phases and adopting the normal or inverted hierarchical scheme of neutrino masses. Taking also m_{iD} as input parameters we can construct the complex symmetric matrix

$$W = -d_D^{-1} \bar{m}_\nu d_D^{-1} = U^c d_{\nu^c} U^{cT} \quad (57)$$

– see Eq. (52) – from which we can extract d_{ν^c} as follows:

$$d_{\nu^c}^{-2} = U^{c\dagger} W W^\dagger U^c. \quad (58)$$

Note that $W W^\dagger$ is a 3×3 complex, hermitian matrix and can be diagonalized following the algorithm described in Ref. [48]. Having determined the elements of U^c and the $M_{i\nu^c}$'s we can compute m_D through Eq. (49) and the ε_i 's through Eq. (44a).

V. CONSTRAINING THE MODEL PARAMETERS

We exhibit the constraints that we impose on our cosmological set-up in Sec. V A, and delineate the allowed parameter space of our model in Sec. V B.

A. IMPOSED CONSTRAINTS

The parameters of our model can be restricted once we impose the following requirements:

1. According to the inflationary paradigm, the horizon and flatness problems of the standard Big Bang cosmology can be successfully resolved provided that \hat{N}_* defined by Eq. (30) takes a certain value, which depends on the details of the cosmological scenario. Employing standard methods [7], we can easily derive the required \hat{N}_* for our model, consistent with the fact that the PQ oscillatory system remains subdominant during the post-inflationary era. Namely we obtain

$$\hat{N}_* \simeq 22.5 + 2 \ln \frac{V_{\text{CI}}(\sigma_*)^{1/4}}{1 \text{ GeV}} - \frac{4}{3} \ln \frac{V_{\text{CI}}(\sigma_f)^{1/4}}{1 \text{ GeV}} + \frac{1}{3} \ln \frac{T_{\text{rh}}}{1 \text{ GeV}} + \frac{1}{2} \ln \frac{f(\sigma_f)}{f(\sigma_*)}. \quad (59)$$

2. The inflationary observables derived in Sec. III B are to be consistent with the fitting [16] of the WMAP7, BAO and H_0 data. As usual, we adopt the central value of $\Delta_{\mathcal{R}}$, whereas we allow the remaining quantities to vary within the 95% confidence level (c.l.) ranges. Namely,

$$\Delta_{\mathcal{R}} \simeq 4.93 \cdot 10^{-5}, \quad (60a)$$

$$0.944 \leq n_s \leq 0.992, \quad (60b)$$

$$-0.062 \leq \alpha_s \leq 0.018, \quad (60c)$$

$$r < 0.24. \quad (60d)$$

3. For the realization of nMCI, we assume that $c_{\mathcal{R}}$ takes relatively large values – see e.g. Eq. (16a). This assumption may [5, 49] jeopardize the validity of the classical approximation, on which the analysis of the inflationary behavior is based. To avoid this inconsistency – which is rather questionable [10, 49] though – we have to check the hierarchy between the ultraviolet cut-off scale [7], $\Lambda = m_{\text{P}}/c_{\mathcal{R}}$, of the effective theory and the inflationary scale, which is represented by $\hat{V}_{\text{CI}}(\sigma_*)^{1/4}$ or, less restrictively, by the corresponding Hubble parameter, $\hat{H}_* = \hat{V}_{\text{CI}}(\sigma_*)^{1/2}/\sqrt{3}m_{\text{P}}$. In particular, the validity of the effective theory implies [49]

$$(a) \quad \hat{V}_{\text{CI}}(\sigma_*)^{1/4} \leq \Lambda \quad \text{or} \quad (b) \quad \hat{H}_* \leq \Lambda. \quad (61)$$

4. To ensure that the inflaton decay according to the lagrangian part of Eq. (39) is kinematically allowed we have to impose the constraint – see Table III:

$$m_{\text{I}} \geq 2m_{\text{PQ}} + M_{i\nu^c} \Rightarrow 2m_{\text{PQ}} + M_{i\nu^c} \lesssim 10^{13} \text{ GeV}, \quad (62)$$

where we make use of Eq. (36). This requirement can be easily satisfied by constraining λ_a and $\lambda_{i\nu^c}$ to values lower than the perturbative limit. As the inequality in Eq. (62) gets strengthened, the accuracy of Eq. (40) where masses of the decay products are neglected, increases.

5. From the solar, atmospheric, accelerator and reactor neutrino experiments we take into account the following inputs [26] – see also Ref. [27] – on the neutrino mass-squared

differences:

$$\Delta m_{21}^2 = (7.59_{-0.18}^{+0.2}) \cdot 10^{-3} \text{ eV}^2, \quad (63a)$$

$$\Delta m_{31}^2 = (2.5_{-0.16}^{+0.09} [-2.4_{-0.09}^{+0.08}]) \cdot 10^{-3} \text{ eV}^2, \quad (63b)$$

on the mixing angles:

$$\sin^2 \theta_{12} = 0.312_{-0.015}^{+0.017}, \quad (63c)$$

$$\sin^2 \theta_{13} = 0.013_{-0.005}^{+0.007} [0.016_{-0.006}^{+0.008}], \quad (63d)$$

$$\sin^2 \theta_{23} = 0.52_{-0.07}^{+0.06} [0.52 \pm 0.06], \quad (63e)$$

and on the CP-violating Dirac phase:

$$\delta = - (0.61_{-0.65}^{+0.75} [0.41_{-0.7}^{+0.65}]) \pi \quad (63f)$$

for normal [inverted] neutrino mass hierarchy. In particular, $m_{i\nu}$'s can be determined via the relations

$$m_{2\nu} = \sqrt{m_{1\nu}^2 + \Delta m_{21}^2} \quad (64a)$$

and

$$m_{3\nu} = \sqrt{m_{1\nu}^2 + \Delta m_{31}^2} \quad (64b)$$

for normally ordered (NO) $m_{i\nu}$'s or

$$m_{1\nu} = \sqrt{m_{3\nu}^2 + |\Delta m_{31}^2|} \quad (64c)$$

for invertedly ordered (IO) $m_{i\nu}$'s. The sum of $m_{i\nu}$'s can be bounded from above by the WMAP7 data [16]

$$\sum_i m_{i\nu} \leq 0.58 \text{ eV} \quad (65)$$

at 95% c.l. This is more restrictive than the 95% c.l. upper bound arising from the effective electron neutrino mass in β -decay [50]:

$$m_{\beta} := |\sum_i U_{1i\nu}^2 m_{i\nu}| \leq 2.3 \text{ eV}. \quad (66)$$

However, in the future, the KATRIN experiment [51] expects to reach the sensitivity of $m_{\beta} \simeq 0.2 \text{ eV}$ at 90% c.l.

6. The interpretation of BAU through nTL dictates [16] at 95% c.l.

$$Y_B = (8.74 \pm 0.42) \cdot 10^{-11}. \quad (67)$$

7. In order to avoid spoiling the success of the BBN, an upper bound on $Y_{\tilde{G}}$ is to be imposed depending on the \tilde{G} mass, $m_{\tilde{G}}$, and the dominant \tilde{G} decay mode. For the conservative case where \tilde{G} decays with a tiny hadronic branching ratio, we have [31]

$$Y_{\tilde{G}} \lesssim \begin{cases} 10^{-14} \\ 2.5 \cdot 10^{-14} \\ 4.3 \cdot 10^{-14} \\ 10^{-13} \end{cases} \quad \text{for } m_{\tilde{G}} \simeq \begin{cases} 0.69 \text{ TeV} \\ 5 \text{ TeV} \\ 8 \text{ TeV} \\ 10.6 \text{ TeV}. \end{cases} \quad (68)$$

As we see below, this bound is achievable within our model model only for $m_{\tilde{G}} \gtrsim 8 \text{ TeV}$. The bound above may be somehow relaxed in the case of a stable \tilde{G} .

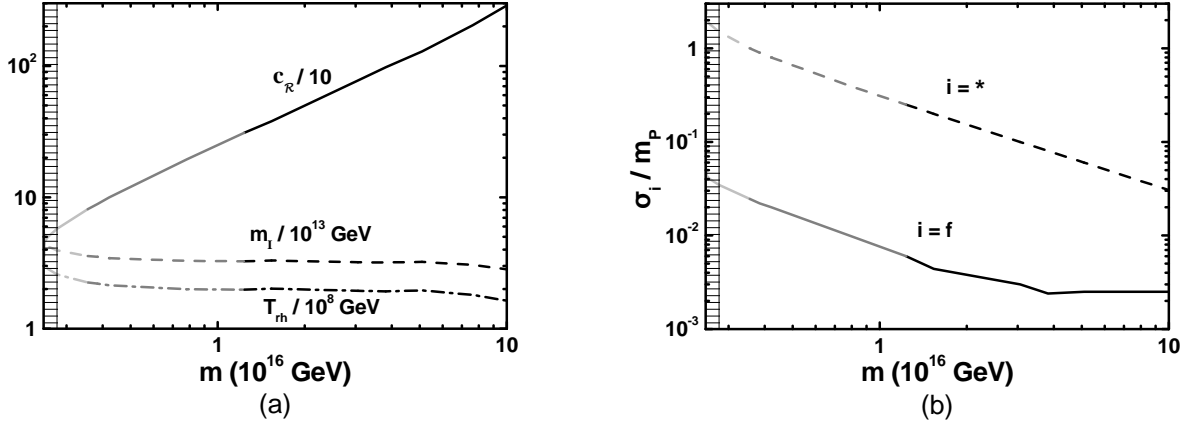


FIG. 1: The allowed by Eqs. (59), (60a), (60b) and (61b) values of $c_{\mathcal{R}}$ (solid line), m_I – given by Eq. (36) – (dashed line) and T_{rh} – given by Eq. (41) – (dot-dashed line) [σ_f (solid line) and σ_* (dashed line)] versus m (a) [(b)] for $\lambda_1 \ll \lambda_2 = \lambda_3 = 0.5$. The light gray and gray segments denote values of the various quantities satisfying Eq. (61a) too, whereas along the light gray segments we obtain $\sigma_* \geq m_P$. Values of the parameters to the right of the lined region correspond to n_s 's lying within its 68% c.l. observationally favored region.

B. RESULTS

As can be easily seen from the relevant expressions in Secs. II and IV B, our cosmological set-up depends on the following independent parameters:

$$m, \lambda_a, \lambda_\mu, k_{\bar{P}}, \lambda_i, f_a, n, m_{\ell\nu}, m_{iD}, \varphi_1 \text{ and } \varphi_2,$$

where $m_{\ell\nu}$ is the low scale mass of the lightest of ν_i 's and can be identified with $m_{1\nu}$ [$m_{3\nu}$] for NO [IO] neutrino mass spectrum. We do not consider $c_{\mathcal{R}}$ and $\lambda_{i\nu^c}$ as independent parameters since $c_{\mathcal{R}}$ is related to m via Eq. (33) while $\lambda_{i\nu^c}$ can be derived from the last six parameters above which affect exclusively the Y_L calculation and can be constrained through the requirements 5 and 6 of Sec. V A. Note that the $\lambda_{i\nu^c}$'s can be replaced by $M_{i\nu^c}$'s given in Table III keeping in mind that perturbativity requires $\lambda_{i\nu^c} \leq \sqrt{4\pi}$ or $M_{i\nu^c} \leq 3.5f_a$. Recall also that V_{rc} in Eq. (24) is independent of $\lambda_{D_a}, \lambda_{h_a} \gg \lambda_a$ and depends only on n , which is set equal to 5 for definiteness. To facilitate the realization of see-saw mechanism, we take $f_a = 10^{12}$ GeV. This choice makes also possible the generation of the μ term of MSSM through the PQ symmetry breaking, since $\mu \sim 1$ TeV is obtained for $\lambda_\mu = 0.01$, whereas lower f_a 's dictate larger λ_μ 's. Moreover, our computation reveals that ε_1 in Eq. (44a) is mostly smaller than ε_2 and ε_3 . Therefore, fulfilling the baryogenesis criterion enforces us to consider $\text{Br}_1 \ll \text{Br}_{2,3}$ or $\lambda_1 \ll \lambda_{2,3} \sim 0.1$. Since ε_2 and ε_3 are of the same order of magnitude, the resulting Y_B does not depend crucially on λ_2/λ_3 . Therefore we believe that $\lambda_2 = \lambda_3 = 0.5$ is a representative choice – e.g., we explicitly checked that the option $\lambda_2 = 0.1$ and $\lambda_3 = 0.9$ or $\lambda_2 = 0.9$ and $\lambda_3 = 0.1$ lead to similar results. Finally, our results are independent of λ_a and $k_{\bar{P}}$ provided Eq. (62) is fulfilled and the positivity of $m_{\bar{P}}^2$ – see Table II – is ensured, respectively. To facilitate the achievement of these objective, we get $\lambda_a = 0.01$ and $k_{\bar{P}} = 1$.

Summarizing, we set throughout our calculation:

$$k_{\bar{P}} = 1, \lambda_1 \leq 0.01, \lambda_2 = \lambda_3 = 0.5, n = 5, \quad (69a)$$

$$\lambda_\mu = \lambda_a = 0.01 \text{ and } f_a = 10^{12} \text{ GeV}. \quad (69b)$$

The selected values for the above quantities give us a wide and natural allowed region for the remaining fundamental parameters of our model, as we show below concentrating separately in the inflationary period (Sec. V B 1) and in the stage of nTL (Sec. V B 2).

1. THE STAGE OF NON-MINIMAL INFLATION

For nMCI, we use as input parameters in our numerical code σ_*, m and $c_{\mathcal{R}}$. For every chosen $c_{\mathcal{R}} \geq 1$ we restrict m and σ_* so that the conditions in Eq. (59) – with T_{rh} evaluated consistently using Eq. (41) – and (60a) are satisfied. Let us remark that, in our numerical calculations, we use the complete formulae for \hat{V}_{CI} – see Eq. (25) –, \hat{N}_* , the slow-roll parameters and $\Delta_{\mathcal{R}}$ in Eqs. (30), (26a), (26b), (32) and not the approximate relations listed in Sec. III B for the sake of presentation.

Our results are displayed in Fig. 1, where we draw the allowed values of $c_{\mathcal{R}}$ (solid line) m_I (dashed line) and T_{rh} (dot-dashed line) [σ_f (solid line) and σ_* (dashed line)] versus m – see Fig. 1-(a) [Fig. 1-(b)]. The constraint of Eq. (61b) is satisfied along the various curves whereas Eq. (61a) is valid only along the gray and light gray segments of these. Along the light gray segments, though, we obtain $\sigma_* \geq m_P$. The lower bound on m is derived from the saturation of the upper bound of inequality in Eq. (60b) whereas the upper bound comes from the fact that the enhanced resulting m 's destabilize the inflationary path through the radiative corrections in Eq. (25) – see Eq. (24). Indeed, V_{rc} starts to influence the inflationary dynamics for $m \geq 1.5 \cdot 10^{16}$ GeV, and consequently, the variation of σ_f as a function of $c_{\mathcal{R}}$ or m – drawn in

Fig. 1-(b) – deviates from the behavior described in Eq. (28). On the contrary the variations of σ_* follows Eq. (31).

In all, we obtain

$$45 \lesssim c_{\mathcal{R}} \lesssim 2950 \quad \text{and} \quad 2.5 \lesssim \frac{m}{10^{15} \text{ GeV}} \lesssim 102 \quad (70)$$

for $\hat{N}_* \simeq 54.5$. From Fig. 1-(a), we observe that m depends on $c_{\mathcal{R}}$ almost linearly whereas m_I remains close to 10^{13} GeV as we anticipated in Eqs. (33) and (36), respectively. As a result of the latter effect, T_{rh} given by Eq. (41) remains also almost constant. As m (or $c_{\mathcal{R}}$) decreases below its maximal value in its allowed region in Eq. (70), we obtain

$$0.965 \lesssim n_s \lesssim 0.991, \quad (71a)$$

$$6.5 \lesssim -\alpha_s/10^{-4} \lesssim 12, \quad (71b)$$

$$3.1 \lesssim r/10^{-3} \lesssim 7.3. \quad (71c)$$

Clearly, the predicted n_s , α_s and r can lie within the allowed ranges given in Eqs. (60b), (60c) and (60d) respectively. In particular, values of the various parameters plotted in Fig. 1, which lie to the right of the lined regions correspond to $n_s \simeq (0.965 - 0.98)$. This result is consistent with the 68% c.l. observationally favored region – see Eq. (60b). It is notable, however, that n_s increases impressively for $\sigma_*/m_{\text{P}} > \sqrt{6}$, contrary to the situation in models of nMCI with quadratic coupling to \mathcal{R} where n_s remains constantly close to its central observational favored value in Eq. (60b) – cf. Ref. [11].

As regards the \tilde{G} abundance, employing Eq. (43), we find

$$3.5 \lesssim Y_{\tilde{G}}/10^{-14} \lesssim 8.4 \quad (72)$$

as m varies within its allowed range in Eq. (70). Comparing this result with the limits of Eq. (68), we infer that our model is consistent with the relevant restriction for $m_{\tilde{G}} \gtrsim (8 - 10)$ TeV.

2. THE STAGE OF NON-THERMAL LEPTOGENESIS

As we show above, the stage of nMCI predicts almost constant values of m_I and T_{rh} – recall that we consider λ_i 's of the order of 0.1. In other words, the post-inflationary evolution in our set-up is largely independent of the precise value of m in the range of Eq. (70). As a consequence, Y_B calculated by Eq. (42) does not vary with m , contrary to the naive expectations. Just for definiteness we take throughout this section $m = 4.2 \cdot 10^{15}$ GeV which corresponds to $c_{\mathcal{R}} = 100$, $n_s = 0.969$, $m_I = 3.4 \cdot 10^{13}$ GeV and $T_{\text{rh}} = 2.1 \cdot 10^8$ GeV ($Y_{\tilde{G}} \simeq 4 \cdot 10^{-14}$) – recall that we use $\lambda_1 \ll \lambda_2 = \lambda_3 = 0.5$.

On the contrary, Y_B in our approach depends crucially on the low energy parameters related to the neutrino physics. In our numerical program, for a given neutrino mass scheme, we take as input parameters: $m_{\ell\nu}, \varphi_1, \varphi_2$ and the best-fit values of the neutrino parameters listed in the paragraph 5 of Sec. V A. We then find the *renormalization group* (RG) evolved values of these parameters at the scale of nTL, Λ_L , which is taken to be $\Lambda_L = m_I$, integrating numerically the

TABLE IV: Parameters yielding the correct BAU for various neutrino mass schemes.

Parameters	Cases						
	A	B	C	D	E	F	G
	Normal Hierarchy		Degenerate Masses			Inverted Hierarchy	
Low Scale Parameters							
$m_{1\nu}/0.1\text{ eV}$	0.05	0.1	0.5	1.	0.7	0.5	0.49
$m_{2\nu}/0.1\text{ eV}$	0.1	0.13	0.51	1.0	0.705	0.51	0.5
$m_{3\nu}/0.1\text{ eV}$	0.5	0.51	0.71	1.12	0.5	0.1	0.05
$\sum_i m_{i\nu}/0.1\text{ eV}$	0.65	0.74	1.7	3.1	1.9	1.1	1
$m_\beta/0.1\text{ eV}$	$8\cdot 10^{-3}$	0.013	0.19	0.46	0.3	0.42	0.44
φ_1	π	π	0	$\pi/4$	$\pi/4$	$\pi/4$	$\pi/4$
φ_2	0	0	$5\pi/6$	π	π	$\pi/2$	$\pi/4$
Leptogenesis-Scale Parameters							
m_{1D}/GeV	2	2.5	4	8	9	6	5
m_{2D}/GeV	3	3.49	5	9.3	6	3	1
m_{3D}/GeV	6.7	4	8	11	4.7	2	2.1
$M_{1\nu^c}/10^{11}\text{ GeV}$	2.5	2.4	3.3	6.5	4.6	1	0.3
$M_{2\nu^c}/10^{11}\text{ GeV}$	11	7.3	5.2	8.13	4.9	5.56	4.3
$M_{3\nu^c}/10^{11}\text{ GeV}$	17	7.6	6	8.36	8.6	6.7	5.1
$\Delta_{iji}/10^4$	4.5	0.5	0.96	0.05	0.3	1.6	1.4
$\Delta_{ijj}/10^4$	1.6	0.54	0.48	0.04	0.3	1.2	3.5
(with $i = 2$ and $j = 3$ except for case E where $i = 1$ and $j = 2$)							
Resulting B -Yield							
$10^{11}Y_B^0$	8.3	7.4	6.3	3.3	7.2	9.3	4.8
$10^{11}Y_B$	8.7	8.85	8.98	8.4	8.9	8.96	8.95

complete expressions of the RG equations – given in Ref. [52] – for $m_{i\nu}, \theta_{ij}, \delta, \varphi_1$ and φ_2 . In doing this, we consider the MSSM with $\tan \beta \simeq 50$ (favored by the preliminary LHC results [53]) as an effective theory between Λ_L and a SUSY-breaking scale, $M_{\text{SUSY}} = 1.5$ TeV. Below M_{SUSY} the running of the various parameters is realized considering the particle content of SM with a mass of about 120 GeV for the light Higgs. Following the procedure described in Sec. IV B, we evaluate $M_{i\nu^c}$ at Λ_L taking m_{iD} as free parameters. In our approach we do not consider the running of m_{iD} and $M_{i\nu^c}$ and therefore we give their values at Λ_L .

We start the exposition of our results arranging in Table IV some representative values of the parameters leading to the correct BAU for normally hierarchical (cases A and B), degenerate (cases C, D and E) and invertedly hierarchical (cases F and G) neutrino masses. For comparison we display the B -yield with (Y_B) or without (Y_B^0) taking into account the RG effects. We observe that the two results differ appreciably especially in the cases with degenerate or IO $m_{i\nu}$'s. As it is evident from the m_{iD} 's chosen, our model is not compatible with any GUT-inspired pattern of large hierarchy be-

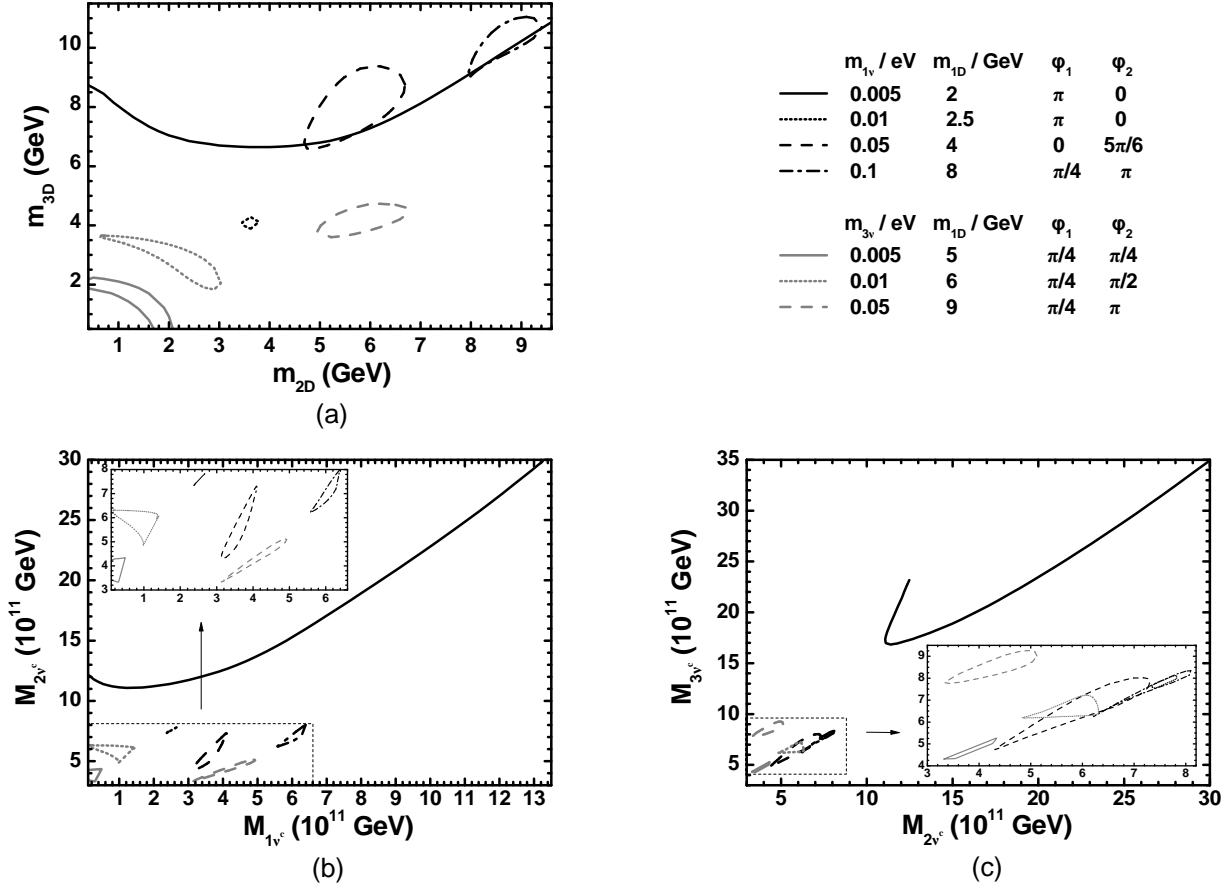


FIG. 2: Contours in the $m_{2D} - m_{3D}$ (a) $M_{1\nu^c} - M_{2\nu^c}$ (b) and $M_{2\nu^c} - M_{3\nu^c}$ (c) plane yielding the central Y_B in Eq. (67), for various $(m_{\ell\nu}, m_{1D}, \phi_1, \phi_2)$'s indicated next to the graph (a) and NO [IO] $m_{i\nu}$'s (black [gray] lines).

tween the m_{iD} 's. In particular, we need $m_{1D} < m_{2D} < m_{3D}$ [$m_{3D} < m_{2D} < m_{1D}$] for NO [IO] $m_{i\nu}$'s (cases A, B, C and D [cases E, F and G]).

From Table IV we also notice that the achievement of Y_B within the range of Eq. (67) dictates mostly proximity between two of the $M_{i\nu^c}$'s. Indeed, except for the case A, we obtain $M_{2\nu^c}/M_{1\nu^c} \simeq 1.06$ in case E and $M_{3\nu^c}/M_{2\nu^c} < 1.2$ in the residual cases. However, it is clear from the displayed Δ_{iji} 's and Δ_{ijj} 's (with $i = 2$ and $j = 3$ for all the cases besides case E where $i = 1$ and $j = 2$) that in our framework the conditions of Eq. (45) are comfortably retained and therefore, our proposal is crucially different from that of resonant leptogenesis [44–46] – it rather resembles that of Ref. [28]. On the other hand, the correctness of Y_B in the case A entails $M_{2\nu^c}$ and $M_{3\nu^c}$ [$\lambda_{2\nu^c}$ and $\lambda_{3\nu^c}$] roughly larger than 10^{12} GeV [unity]. In all cases the current limit of Eq. (65) is safely met – the case D approaches it –, while m_β turns out to be well below the projected sensitivity of KATRIN [51].

To highlight further our conclusions inferred from Table IV, we can fix $m_{\ell\nu}$ ($m_{1\nu}$ for NO $m_{i\nu}$'s or $m_{3\nu}$ for IO $m_{i\nu}$'s) m_{1D} , ϕ_1 and ϕ_2 to their values shown in this table and vary m_{2D} and m_{3D} so that the central value of Eq. (67) is achieved. The resulting contours in the $m_{2D} - m_{3D}$ plane are presented in Fig. 2-(a) – since the range of Eq. (67) is very narrow the

possible variation of the drawn lines is negligible. The resulting values of $M_{j\nu^c}$ are displayed in $M_{1\nu^c} - M_{2\nu^c}$ and $M_{2\nu^c} - M_{3\nu^c}$ plane – see Fig. 2-(b) and Fig. 2-(c) respectively. The conventions adopted for the types and the color of the various lines are also described next to the graphs (a) of Fig. 2. In particular, we use black [gray] lines for NO [IO] $m_{i\nu}$'s. Besides the case with $m_{1\nu} = 0.005$ eV we observe that every curve in all graphs has two branches and not large hierarchies allowed in the sectors of both the m_D 's and M_{ν^c} 's. Note that the black contour for $m_{1\nu} = 0.01$ eV in Fig. 2-(c) is included within the one for $m_{1\nu} = 0.1$ eV and so, it is not quite distinguishable. For $m_{1\nu} = 0.005$ eV, $\lambda_{3\nu^c}$ saturates the perturbation limit. Since we expect that the $\lambda_{i\nu^c}$'s increase [54] due to their RG running from low to higher scale, our results do not jeopardize the validity of the conventional perturbation approach up to the scale Λ_L . In all cases we find $1 \lesssim m_{iD}/\text{GeV} \lesssim 10$.

It is worth emphasizing that, although our mechanism of nTL is connected with the specific inflationary model under consideration, it can have a much wider applicability. It can be realized within other models of inflation with similar inflaton mass and reheat temperature, since it is largely independent of the details of the inflationary phase but restricts mainly the yet unknown parameters of neutrino physics ($m_{i\nu}, m_{iD}, \phi_1, \phi_2$).

VI. CONCLUSIONS

We investigated a novel inflationary scenario in which the inflaton field appears in a bilinear superpotential term and in a linear holomorphic function included in a logarithmic Kähler potential. The latter function can be interpreted in JF as a non-minimal coupling to gravity, whose the strength is constrained so as the EF inflationary potential can be flattened enough to support a stage of non-minimal inflation compatible with observations. The inflationary model was embedded in a moderate extension of MSSM augmented by three RH neutrino superfields and three other singlet superfields, which lead to a PQPT tied to renormalizable superpotential terms. The PQPT follows nMCI and resolves the strong CP and the μ problems of MSSM and also provides RH neutrinos with masses lower than about 10^{12} GeV. The possible catastrophic production of domain walls can be eluded by the introduction of extra matter superfields which can be chosen so that the MSSM gauge coupling constant unification is not disturbed. For \tilde{G} masses larger than 8 TeV, observationally safe reheating of the universe with $T_{\text{rh}} \simeq 10^8$ GeV can be accomplished by a three-body decay of the inflaton. The subsequent out-of-

equilibrium decays of the produced RH sneutrinos can generate the required by the observations BAU consistently with the present low energy neutrino data, provided that the Dirac neutrino masses are constrained in the range (1 – 10) GeV for all the light neutrino mass schemes. It is gratifying that the degeneracy of the masses of the RH (s)neutrinos required by the mechanism of nTL in our model is low enough compared with their decay widths, so that perturbative calculation remains safely valid. Finally, we briefly discussed scenarios in which the potential axino and saxion overproduction problems can be avoided.

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REFERENCES

- [1] K. Nozari and S.D. Sadatian, *Mod. Phys. Lett. A* **23**, 2933 (2008) [arXiv:0710.0058]; K. Nozari and S. Shafizadeh, *Phys. Scripta* **82**, 015901 (2010) [arXiv:1006.1027].
- [2] J.L. Cervantes-Cota and H. Dehnen, *Nucl. Phys. B* **442**, 391 (1995) [astro-ph/9505069]; F.L. Bezrukov and M. Shaposhnikov, *Phys. Lett. B* **659**, 703 (2008) [arXiv:0710.3755]; A. Barvinsky *et al.*, *J. Cosmol. Astropart. Phys.* **11**, 021 (2008) [arXiv:0809.2104]; A. De Simone, M.P. Hertzberg and F. Wilczek, *Phys. Lett. B* **678**, 1 (2009) [arXiv:0812.4946].
- [3] R. Fakir and W.G. Unruh, *Phys. Rev. D* **41**, 1783 (1990); E. Komatsu and T. Futamase, *Phys. Rev. D* **59**, 064029 (1999) [astro-ph/9901127]; S. Tsujikawa and B. Gumjudpai, *Phys. Rev. D* **69**, 123523 (2004) [astro-ph/0402185].
- [4] T.E. Clark, B. Liu, S.T. Love and T. ter Veldhuis, *Phys. Rev. D* **80**, 075019 (2009) [arXiv:0906.5595]; R.N. Lerner and J. McDonald, *Phys. Rev. D* **80**, 123507 (2009) [arXiv:0909.0520].
- [5] R.N. Lerner and J. McDonald, *Phys. Rev. D* **82**, 103525 (2010) [arXiv:1005.2978]; G.F. Giudice and H.M. Lee, *Phys. Lett. B* **694**, 294 (2011) [arXiv:1010.1417]; F. Bauer and D.A. Demir, *Phys. Lett. B* **698**, 425 (2011) [arXiv:1012.2900].
- [6] S.C. Park and S. Yamaguchi, *J. Cosmol. Astropart. Phys.* **08**, 009 (2008) [arXiv:0801.1722]; S. Koh and M. Minamitsuji, *Phys. Rev. D* **83**, 046009 (2011) [arXiv:1011.4655]; A. De Felice, S. Tsujikawa, J. Elliston and R. Tavakol, *J. Cosmol. Astropart. Phys.* **08**, 021 (2011) [arXiv:1105.4685].
- [7] C. Pallis, *Phys. Lett. B* **692**, 287 (2010) [arXiv:1002.4765].
- [8] N. Okada, M.U. Rehman and Q. Shafi, *Phys. Rev. D* **82**, 043502 (2010) [arXiv:1005.5161]; N. Okada, M.U. Rehman and Q. Shafi, *Phys. Lett. B* **701**, 520 (2011) [arXiv:1102.4747].
- [9] M.B. Einhorn and D.R.T. Jones, *J. High Energy Phys.* **03**, 026 (2010) [arXiv:0912.2718]; S. Ferrara *et al.*, *Phys. Rev. D* **82**, 045003 (2010) [arXiv:1004.0712]; H.M. Lee, *J. Cosmol. Astropart. Phys.* **08**, 003 (2010) [arXiv:1005.2735]; S. Ferrara *et al.*, *Phys. Rev. D* **83**, 025008 (2011) [arXiv:1008.2942].
- [10] R. Kallosh and A. Linde, *J. Cosmol. Astropart. Phys.* **11**, 011 (2010) [arXiv:1008.3375]; A. Linde, M. Noorbala and A. Westphal, *J. Cosmol. Astropart. Phys.* **03**, 013 (2011) [arXiv:1101.2652].
- [11] C. Pallis and N. Toumbas, *J. Cosmol. Astropart. Phys.* **02**, 019 (2011) [arXiv:1101.0325].
- [12] C. Pallis and N. Toumbas, *J. Cosmol. Astropart. Phys.* **12**, 002 (2011) [arXiv:1108.1771].
- [13] Y.G. Kim, H.M. Lee, W.-I. Park, *J. High Energy Phys.* **08**, 126 (2011) [arXiv:1107.1113]; M. Arai, S. Kawai and N. Okada, *Phys. Rev. D* **84**, 123515 (2011) [arXiv:1107.4767]; M. Arai, S. Kawai and N. Okada, arXiv:1112.2391.
- [14] M. Kawasaki, M. Yamaguchi and T. Yanagida, *Phys. Rev. Lett.* **85**, 3572 (2000) [hep-ph/0004243]; M. Kawasaki, M. Yamaguchi and T. Yanagida, *Phys. Rev. D* **63**, 103514 (2001) [hep-ph/0011104]; S. Antusch, M. Bastero-Gil, K. Dutta, S.F. King and P.M. Kostka, *Phys. Lett. B* **679**, 428 (2009) [arXiv:0905.0905]; F. Takahashi, *Phys. Lett. B* **693**, 140 (2010) [arXiv:1006.2801].
- [15] P. Brax and J. Martin, *Phys. Rev. D* **72**, 023518 (2005) [hep-th/0504168]; I. Ben-Dayana and M.B. Einhorn, *J. Cosmol. Astropart. Phys.* **12**, 002 (2010) [arXiv:1009.2276]; K. Nakayama and F. Takahashi, *J. Cosmol. Astropart. Phys.* **02**, 010 (2011) [arXiv:1008.4457].
- [16] E. Komatsu *et al.* [WMAP Collaboration], *Astrophys. J. Suppl.* **192**, 18 (2011) [arXiv:1001.4538]; <http://lambda.gsfc.nasa.gov/product/map/dr2/parameters.cfm>.
- [17] R. Peccei and H. Quinn, *Phys. Rev. Lett.* **38**, 1440 (1977); S. Weinberg, *ibid.* **40**, 223 (1978); F. Wilczek, *ibid.* **40**, 279 (1978).
- [18] H. Baer, A.D. Box and H. Summy, *J. High Energy Phys.*

- 10**, 023 (2010) [arXiv:1005.2215]; H. Baer, A. Lessa and W. Sreethawong, *J. Cosmol. Astropart. Phys.* **01**, 036 (2012) [arXiv:1110.2491].
- [19] M. Kawasaki, N. Kitajima and K. Nakayama, *Phys. Rev. D* **83**, 123521 (2011) [arXiv:1104.1262]; E.J. Chun, *Phys. Rev. D* **84**, 043509 (2011) [arXiv:1104.2219]; K.-Y. Choi, L. Covi, J.E. Kim and L. Roszkowski, *J. High Energy Phys.* **04**, 106 (2012) [arXiv:1108.2282]; K.J. Bae, E.J. Chun and S.H. Im, *J. Cosmol. Astropart. Phys.* **03**, 013 (2012) [arXiv:1111.5962].
- [20] J.E. Kim, *Phys. Lett. B* **136**, 378 (1984); J.E. Kim and H.P. Nilles, *Phys. Lett. B* **138**, 150 (1984); T. Goto and M. Yamaguchi, *Phys. Lett. B* **276**, 103 (1992).
- [21] P. Minkowski, *Phys. Lett. B* **67**, 421 (1977); T. Yanagida, Proceedings of the Workshop on the Unified Theory and the Baryon Number in the Universe (O. Sawada and A. Sugamoto, eds.), KEK, Tsukuba, Japan, 1979, p. 95; M. Gell-Mann, P. Ramond, and R. Slansky, *Supergravity* (P. van Nieuwenhuizen *et al.* eds.), North Holland, Amsterdam, 1979, p. 315; S.L. Glashow, Proceedings of the 1979 Cargese Summer Institute on Quarks and Leptons (M. Levy *et al.* eds.), Plenum Press, New York, 1980, p. 687; R.N. Mohapatra and G. Senjanović, *Phys. Rev. Lett.* **44**, 912 (1980).
- [22] P. Sikivie, *Phys. Rev. Lett.* **48**, 1156 (1982).
- [23] G. Lazarides and Q. Shafi, *Phys. Lett. B* **115**, 21 (1982); H. Georgi and M.B. Wise, *ibid.* **116**, 123 (1982).
- [24] G. Lazarides and Q. Shafi, *Phys. Lett. B* **489**, 194 (2000) [hep-ph/0006202]; G. Lazarides and C. Pallis, *Phys. Rev. D* **82**, 063535 (2010) [arXiv:1007.1558].
- [25] G. Lazarides and Q. Shafi, *Phys. Lett. B* **258**, 305 (1991); K. Kumekawa, T. Moroi and T. Yanagida, *Prog. Theor. Phys.* **92**, 437 (1994) [hep-ph/9405337]; G. Lazarides, R.K. Schaefer and Q. Shafi, *Phys. Rev. D* **56**, 1324 (1997) [hep-ph/9608256]; G. Lazarides, hep-ph/9905450.
- [26] T. Schwetz, M. Tortola, J.W.F. Valle, *New J. Phys.* **13**, 109401 (2011) [arXiv:1108.1376].
- [27] G.L. Fogli, E. Lisi, A. Marrone, A. Palazzo and A.M. Rotunno, *Phys. Rev. D* **84**, 053007 (2011) [arXiv:1106.6028].
- [28] T. Asaka, K. Hamaguchi, M. Kawasaki and T. Yanagida, *Phys. Rev. D* **61**, 083512 (2000) [hep-ph/9907559]; M. Fujii, K. Hamaguchi and T. Yanagida, *Phys. Rev. D* **65**, 115012 (2002) [hep-ph/0202210].
- [29] M. Yu. Khlopov and A.D. Linde, *Phys. Lett. B* **138**, 265 (1984); J. Ellis, J.E. Kim and D.V. Nanopoulos, *Phys. Lett. B* **145**, 181 (1984).
- [30] M. Bolz, A. Brandenburg and W. Buchmüller, *Nucl. Phys. B* **606**, 518 (2001); M. Bolz, A. Brandenburg and W. Buchmüller, *Nucl. Phys. B* **790**, 336 (2008) (E) [hep-ph/0012052]; J. Pradler and F.D. Steffen, *Phys. Rev. D* **75**, 023509 (2007) [hep-ph/0608344].
- [31] M. Kawasaki, K. Kohri and T. Moroi, *Phys. Lett. B* **625**, 7 (2005) [astro-ph/0402490]; M. Kawasaki, K. Kohri and T. Moroi, *Phys. Rev. D* **71**, 083502 (2005) [astro-ph/0408426]; R.H. Cyburt *et al.*, *Phys. Rev. D* **67**, 103521 (2003) [astro-ph/0211258]; J.R. Ellis, K.A. Olive and E. Vangioni, *Phys. Lett. B* **619**, 30 (2005) [astro-ph/0503023].
- [32] See e.g. M. Peskin, arXiv:0801.1928.
- [33] S.R. Coleman and E.J. Weinberg, *Phys. Rev. D* **7**, 1888 (1973).
- [34] D.H. Lyth and A. Riotto, *Phys. Rept.* **314**, 1 (1999) [hep-ph/9807278]; A. Mazumdar and J. Rocher, *Phys. Rept.* **497**, 85 (2011) [arXiv:1001.0993].
- [35] D.S. Salopek, J.R. Bond and J.M. Bardeen, *Phys. Rev. D* **40**, 1753 (1989); F. Bauer and D.A. Demir, *Phys. Lett. B* **665**, 222 (2008) [arXiv:0803.2664]; N. Makino and M. Sasaki, *Prog. Theor. Phys.* **86**, 103 (1991); R. Fakir, S. Habib and W. Unruh, *Astrophys. J.* **394**, 396 (1992); D.I. Kaiser, *Phys. Rev. D* **52**, 4295 (1995) [astro-ph/9408044]; T. Chiba and M. Yamaguchi, *J. Cosmol. Astropart. Phys.* **10**, 021 (2008) [arXiv:0807.4965].
- [36] L. Kofman, A.D. Linde and A.A. Starobinsky, *Phys. Rev. Lett.* **73**, 3195 (1994) [hep-th/9405187]; L. Kofman, A.D. Linde and A.A. Starobinsky, *Phys. Rev. D* **56**, 3258 (1997) [hep-ph/9704452].
- [37] J. Garcia-Bellido, D.G. Figueroa and J. Rubio, *Phys. Rev. D* **79**, 063531 (2009) [arXiv:0812.4624]; F. Bezrukov, D. Gorbunov and M. Shaposhnikov, *J. Cosmol. Astropart. Phys.* **06**, 029 (2009) [arXiv:0812.3622].
- [38] G.N. Felder, L. Kofman and A.D. Linde, *Phys. Rev. D* **59**, 123523 (1999) [hep-ph/9812289].
- [39] C. Pallis, *Nucl. Phys. B* **751**, 129 (2006) [hep-ph/0510234].
- [40] K. Hamaguchi, *Phd Thesis* [hep-ph/0212305]; W. Buchmüller, R.D. Peccei and T. Yanagida, *Ann. Rev. Nucl. Part. Sci.* **55**, 311 (2005) [hep-ph/0502169].
- [41] M. Endo, M. Kawasaki, F. Takahashi and T.T. Yanagida, *Phys. Lett. B* **642**, 518 (2006) [hep-ph/0607170]; M. Endo, F. Takahashi and T.T. Yanagida, *Phys. Rev. D* **76**, 083509 (2007) [arXiv:0706.0986].
- [42] H. Baer, S. Kraml, A. Lessa and S. Sekmen, *J. Cosmol. Astropart. Phys.* **04**, 039 (2011) [arXiv:1012.3769].
- [43] M. Kawasaki, K. Nakayama and M. Senami, *J. Cosmol. Astropart. Phys.* **03**, 009 (2008) [arXiv:0711.3083].
- [44] M. Flanz, E. A. Paschos and U. Sarkar, *Phys. Lett. B* **345**, 248 (1995); Erratum-ibid. **B 382**, 447 (1996) [hep-ph/9411366]; M. Flanz, E.A. Paschos, U. Sarkar and J. Weiss, *Phys. Lett. B* **389**, 693 (1996) [hep-ph/9607310]; L. Covi, E. Roulet and F. Vissani, *Phys. Lett. B* **384**, 169 (1996) [hep-ph/9605319].
- [45] A. Pilaftsis, *Phys. Rev. D* **56**, 5431 (1997) [hep-ph/9707235]; A. Pilaftsis and T.E.J. Underwood, *Nucl. Phys. B* **692**, 303 (2004) [hep-ph/0309342]; F.F. Deppisch and A. Pilaftsis, *Phys. Rev. D* **83**, 076007 (2011) [arXiv:1012.1834].
- [46] W. Buchmüller and M. Plumacher, *Phys. Lett. B* **431**, 354 (1998) [hep-ph/9710460]; A. Anisimov, A. Broncano and M. Plumacher, *Nucl. Phys. B* **737**, 176 (2006) [hep-ph/051124]; M. Garmy, A. Hohenegger, A. Kartavtsev and M. Lindner, *Phys. Rev. D* **81**, 085027 (2010) [arXiv:0911.4122].
- [47] G.C. Branco, R. Gonzalez Felipe, F.R. Joaquim and M.N. Rebelo, *Nucl. Phys. B* **640**, 202 (2002) [hep-ph/0202030]; E.K. Akhmedov, M. Frigerio and A.Y. Smirnov, *J. High Energy Phys.* **09**, 021 (2003) [hep-ph/0305322]; S. Pascoli, S.T. Petcov and W. Rodejohann, *Phys. Rev. D* **68**, 093007 (2003) [hep-ph/0302054]; V.N. Şenoğuz, *Phys. Rev. D* **76**, 013005 (2007) [arXiv:0704.3048].
- [48] J. Kopp, *Int. J. Mod. Phys. C* **19**, 523 (2008) [physics/0610206].
- [49] C.P. Burgess, H.M. Lee and M. Trott, *J. High Energy Phys.* **09**, 103 (2009) [arXiv:0902.4465]; J.L.F. Barbon and J.R. Espinosa, *Phys. Rev. D* **79**, 081302 (2009) [arXiv:0903.0355]; C.P. Burgess, H.M. Lee and M. Trott, *J. High Energy Phys.* **07**, 007 (2010) [arXiv:1002.2730]; R.N. Lerner and J. McDonald, *J. Cosmol. Astropart. Phys.* **04**, 015 (2010) [arXiv:0912.5463]; M.P. Hertzberg, *J. High Energy Phys.* **11**, 023 (2010) [arXiv:1002.2995]; F. Bezrukov *et al.*, *J. High Energy Phys.* **01**, 016 (2011) [arXiv:1008.5157]; M. Atkins and X. Calmet, *Phys. Lett. B* **697**, 37 (2011) [arXiv:1011.4179].
- [50] H.V. Klapdor-Kleingrothaus *et al.*, *Eur. Phys. J. A* **12**, 147 (2001) [hep-ph/0103062]; C.E. Aalseth *et al.* [IGEX Collaboration], *Phys. Rev. D* **65**, 092007 (2002) [hep-ex/0202026];

- J. Bonn *et al.*, *Nucl. Phys. Proc. Suppl.* **91** (2001) 273; V. M. Lobashev *et al.*, *Nucl. Phys. Proc. Suppl.* **91**, 280 (2001).
- [51] A. Osipowicz *et al.* [KATRIN Collaboration], hep-ex/0109033.
- [52] S. Antusch, J. Kersten, M. Lindner and M. Ratz, *Nucl. Phys.* **B674**, 401 (2003) [hep-ph/0305273]; <http://www.ph.tum.de/~mratz/AnalyticFormulae>.
- [53] A. Arbey *et al.* *Phys. Lett. B* **708**, 162 (2012) [arXiv:1112.3028]; M. Kadastik, K. Kannike, A. Racioppi and M. Raidal, *J. High Energy Phys.* **05**, 061 (2012) [arXiv:1112.3647]; I. Gogoladze, Q. Shafi, and C.S. Un, arXiv:1112.2206; O. Buchmueller *et al.*, arXiv:1112.3564; N. Karagiannakis, G. Lazarides and C. Pallis, arXiv:1201.2111.
- [54] H. Murayama, H. Suzuki and T. Yanagida, *Phys. Lett. B* **291**, 418 (1992); K. Choi, E.J. Chun and J.E. Kim, *Phys. Lett. B* **403**, 209 (1997) [hep-ph/9608222]; W.I. Park, *J. High Energy Phys.* **07**, 085 (2010) [arXiv:1004.2326].